

Probing the charged Higgs boson at the LHC in the CP-violating type-II 2HDM

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ABSTRACT: We present a phenomenological study of a CP-violating two-Higgs-doublet Model with type-II Yukawa couplings at the Large Hadron Collider (LHC). In the light of recent LHC data, we focus on the parameter space that survives the current and past experimental constraints as well as theoretical bounds on the model. Once the phenomenological scenario is set, we analyse the scope of the LHC in exploring this model through the discovery of a charged Higgs boson produced in association with a W boson, with the former decaying into the lightest neutral Higgs and a second W state, altogether yielding a $b\bar{b}W^+W^-$ signature, of which we exploit the W^+W^- semileptonic decays.

KEYWORDS: Quantum field theory, Higgs Physics, 2HDM, CP violation

Contents

1	Introduction	1
2	The model	3
2.1	Yukawa and gauge couplings	5
3	Constraining the parameter space	5
3.1	Theoretical constraints	5
3.2	Experimental constraints	6
3.3	Two scenarios	10
3.4	Studies in α space	11
3.4.1	Scenario 1	11
3.4.2	Scenario 2	13
3.4.3	CP-conserving limits	14
3.5	Studies in $\tan\beta$ - M_{H^\pm} space	15
3.5.1	The unitarity constraint	15
3.5.2	The experimental constraints	16
4	Benchmark analysis	17
4.1	Charged and lightest neutral Higgs bosons: BRs	18
4.2	Charged Higgs: single production mechanisms	20
4.3	$pp \rightarrow H^\pm W^\mp$: significance analysis	23
5	Possible future scenarios	29
5.1	Higher and more constrained rates for $gg \rightarrow H_1 \rightarrow \gamma\gamma$	30
5.2	Tightened upper bound on $gg \rightarrow H_{2,3} \rightarrow W^+W^-$ (and ZZ)	32
6	Conclusions	32
A	Appendix. The decoupling limit	34
A.1	$\tan\beta$ of $\mathcal{O}(1)$	34
A.2	Large $\tan\beta$	35

1 Introduction

One of the main tasks of the Large Hadron Collider (LHC) experiments is to study and understand the mechanism of Electroweak Symmetry Breaking (EWSB). Recently, the ATLAS and CMS collaborations released the results of the search for the Higgs bosons with more than 10 fb^{-1} data collected at 7 TeV in 2011 and at 8 TeV in 2012 [1, 2]. Both experiments have recorded an excess of events above the expected background in

different decay channels (mainly $\gamma\gamma$, ZZ and W^+W^-) at a mass near 125 GeV. The excess is compatible with the Standard Model (SM) Higgs boson. Complementary evidence is also provided by the updated combination of the Higgs searches performed by the CDF and D0 collaborations at the Tevatron [3].

Investigating the Higgs mechanism in the framework of the SM constitutes a major effort [4]. However, the minimal choice of Higgs sector is (so far) arbitrary. Even if the existence of a scalar resonance, compatible with an SM-like Higgs, has already been uncovered by the current data, one must take advantage of the unique opportunity to test the phenomenology of more complicated Higgs models.

Much effort has been dedicated over the years to the study of extended Higgs sectors. In this paper we consider the two-Higgs-doublet Model (2HDM) with type-II Yukawa sector. This model is one of the most popular extensions of the Higgs sector due to its strong connection with a tree-level Minimal Supersymmetric Standard Model (MSSM) [5, 6], which is one of the most accredited proposals for solving some theoretical inconsistencies of the SM. As is well known, the Higgs sector of the MSSM is quite well constrained in terms of the number of free parameters on which the masses depend [7–22]. It is possible that the Higgs sector lies in a lower mass range than the superpartners of the SM particles. In this regard, the 2HDM should be explored as an effective low-energy MSSM-like Higgs sector.

While a tree-level MSSM Higgs sector is strictly CP-conserving (no mixing is allowed between the scalar and pseudo-scalar Higgs components), it has been shown that a CP-violating effective Higgs sector could be produced by loop corrections under specific circumstances [23]. Accordingly, in this paper we adopt a bottom-up approach by considering a CP-violating 2HDM with type-II Yukawa couplings.

Due to its complicated Yukawa structure, a CP-violating parameter space must be carefully constrained by theoretical arguments and experimental data. Therefore, our first aim is to provide a detailed analysis of the allowed parameter space in the light of recent LHC results. We will show that very little CP-conserving parameter space survives these data. This exploration of the allowed parameter space has been addressed recently by several authors, from different points of view [24–28]. Where there is overlap, we compare our results with those obtained by these authors.

Regarding phenomenology, the only way to unambiguously probe the existence of a Higgs sector with two doublets arises through the discovery of a charged Higgs boson, since this particle is the hallmark of such a structure of the Higgs sector. Hence, our second aim is to profile a charged Higgs boson in the surviving parameter space via a detailed study of its production cross-section and decay Branching Ratios (BRs).

Then, our third aim is to study the scope of the LHC in discovering a charged Higgs state. In this respect, it is well known that the production of a single charged Higgs state at a hadron collider proceeds in association with either top/bottom quarks or scalar/vector bosons [29, 30]. By taking into account the recent experimental excess observed by ATLAS and CMS [1, 2] and the Tevatron [3] (i.e., corresponding to a light Higgs with a mass of ≈ 125 GeV), we propose a search strategy for a charged Higgs boson produced in association with a W boson and decaying into a $b\bar{b}W$ final state. In particular, we show that an

appropriate choice of the selection cuts would allow the discovery of such a particle despite the considerable $t\bar{t}$ dominated background.

This paper is organised as follows: in section 2 we give an overview of the considered model, in section 3 we analyse the allowed parameter space in the light of both theoretical and experimental constraints, in section 4 we present the main phenomenological results, in section 5 we briefly comment on possible future developments, and in section 6 we present our conclusions. In the appendix we discuss the decoupling limit for the CP-violating type-II 2HDM.

2 The model

We describe here our parametrisation of the 2HDM with type-II Yukawa couplings. The Higgs sector is defined by the presence of two Higgs doublets, with one (Φ_2) field coupled to the u -type quarks, and the other (Φ_1) to the d -type quarks and charged leptons [5].

We take the 2HDM potential to be

$$\begin{aligned} V = & \frac{\lambda_1}{2}(\Phi_1^\dagger\Phi_1)^2 + \frac{\lambda_2}{2}(\Phi_2^\dagger\Phi_2)^2 + \lambda_3(\Phi_1^\dagger\Phi_1)(\Phi_2^\dagger\Phi_2) \\ & + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) + \frac{1}{2} \left[\lambda_5(\Phi_1^\dagger\Phi_2)^2 + \text{h.c.} \right] \\ & - \frac{1}{2} \left\{ m_{11}^2(\Phi_1^\dagger\Phi_1) + \left[m_{12}^2(\Phi_1^\dagger\Phi_2) + \text{h.c.} \right] + m_{22}^2(\Phi_2^\dagger\Phi_2) \right\}. \end{aligned} \quad (2.1)$$

The Z_2 symmetry will be respected by the quartic terms (there are no λ_6 or λ_7 terms), and Flavour-Changing Neutral Currents (FCNCs) are constrained [31].

We parametrise the Higgs fields as

$$\Phi_1 = \begin{pmatrix} -s_\beta H^+ \\ \frac{1}{\sqrt{2}}[v_1 + \eta_1 - i s_\beta \eta_3] \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} c_\beta H^+ \\ \frac{1}{\sqrt{2}}[v_2 + \eta_2 + i c_\beta \eta_3] \end{pmatrix}. \quad (2.2)$$

The real and non-negative Vacuum Expectation Values (VEVs) for the Higgs doublets are $v_1 = v c_\beta$ and $v_2 = v s_\beta$, with $c_\beta = \cos \beta$ and $s_\beta = \sin \beta$, and the ratio defines

$$\tan \beta = \frac{v_2}{v_1}. \quad (2.3)$$

CP violation is allowed, and it is realised by means of the fact that λ_5 and m_{12}^2 are complex numbers. All three neutral states will then mix, with the physical Higgs particles H_i ($i = 1, 2, 3$) related to the weak fields η_j ($j = 1, 2, 3$) of Eq. (2.2) by

$$\begin{pmatrix} H_1 \\ H_2 \\ H_3 \end{pmatrix} = R \begin{pmatrix} \eta_1 \\ \eta_2 \\ \eta_3 \end{pmatrix}. \quad (2.4)$$

In terms of the non-diagonal mass-squared matrix \mathcal{M}^2 , determined from second derivatives of the above potential, we have

$$R\mathcal{M}^2 R^T = \mathcal{M}_{\text{diag}}^2 = \text{diag}(M_1^2, M_2^2, M_3^2). \quad (2.5)$$

The 3×3 mixing matrix R governing the neutral sector will be parametrised in terms of the angles α_1 , α_2 and α_3 as in [32, 33]:

$$R = \begin{pmatrix} c_1 c_2 & s_1 c_2 & s_2 \\ -(c_1 s_2 s_3 + s_1 c_3) & c_1 c_3 - s_1 s_2 s_3 & c_2 s_3 \\ -c_1 s_2 c_3 + s_1 s_3 & -(c_1 s_3 + s_1 s_2 c_3) & c_2 c_3 \end{pmatrix} \quad (2.6)$$

where $c_1 = \cos \alpha_1$, $s_1 = \sin \alpha_1$, etc., and

$$-\frac{\pi}{2} < \alpha_1 \leq \frac{\pi}{2}, \quad -\frac{\pi}{2} < \alpha_2 \leq \frac{\pi}{2}, \quad 0 \leq \alpha_3 \leq \frac{\pi}{2}. \quad (2.7)$$

For these angular ranges, we have $c_i \geq 0$, $s_3 \geq 0$, whereas s_1 and s_2 may be either positive or negative. We will use the terminology “general 2HDM” as a reminder that CP violation is allowed.

With all three masses different, there are three limits of *no* CP-violation, i.e., with two Higgs bosons that are CP-even and one that is odd. In the above notation, the three limits are [34]:

$$\begin{aligned} H_1 \text{ odd: } & \alpha_2 \simeq \pm\pi/2, \alpha_1, \alpha_3 \text{ arbitrary,} \\ H_2 \text{ odd: } & \alpha_2 = 0, \alpha_3 = \pi/2, \alpha_1 \text{ arbitrary,} \\ H_3 \text{ odd: } & \alpha_2 = \alpha_3 = 0, \alpha_1 \text{ arbitrary.} \end{aligned} \quad (2.8)$$

In the general CP-violating case, the neutral sector is conveniently described by these three mixing angles, together with two masses (M_1, M_2), $\tan \beta$ (the ratio between the two Higgs VEVs) and the parameter $\mu^2 = \text{Re } m_{12}^2 / (2 \cos \beta \sin \beta)$. From Eq. (2.5), it follows that

$$(\mathcal{M}^2)_{ij} = \sum_k R_{ki} M_k^2 R_{kj}. \quad (2.9)$$

When CP is violated, both $(\mathcal{M}^2)_{13}$ and $(\mathcal{M}^2)_{23}$ will be non-zero. In fact, they are related by

$$(\mathcal{M}^2)_{13} = \tan \beta (\mathcal{M}^2)_{23}. \quad (2.10)$$

From these two equations, (2.9) and (2.10), we can determine M_3 from M_1, M_2 , the angles $(\alpha_1, \alpha_2, \alpha_3)$ and $\tan \beta$ [32]:

$$M_3^2 = \frac{M_1^2 R_{13}(R_{12} \tan \beta - R_{11}) + M_2^2 R_{23}(R_{22} \tan \beta - R_{21})}{R_{33}(R_{31} - R_{32} \tan \beta)} \quad (2.11)$$

where we impose $M_1 \leq M_2 \leq M_3$.

Providing also M_{H^\pm} and μ^2 , all the λ ’s are consequently determined. Since the left-hand side of (2.9) can be expressed in terms of the parameters of the potential (see, for example, [35]), we can solve these equations and obtain the λ ’s in terms of the rotation matrix, the neutral mass eigenvalues, μ^2 and M_{H^\pm} . The explicit expressions are given in Ref. [34].

The interest in allowing for CP violation lies in the fact that it may be helpful for baryogenesis [36]. Also, from a more pragmatic point of view, it opens up a bigger parameter space, and allows certain couplings to be larger.

2.1 Yukawa and gauge couplings

For the type-II 2HDM, and for the third generation, the neutral-sector Yukawa couplings are (assuming all fields incoming):

$$\begin{aligned} H_j b \bar{b} : & \quad \frac{-ig m_b}{2 m_W} \frac{1}{\cos \beta} [R_{j1} - i\gamma_5 \sin \beta R_{j3}], \\ H_j t \bar{t} : & \quad \frac{-ig m_t}{2 m_W} \frac{1}{\sin \beta} [R_{j2} - i\gamma_5 \cos \beta R_{j3}]. \end{aligned} \quad (2.12)$$

Likewise, the charged-Higgs couplings are [5]

$$\begin{aligned} H^+ b \bar{t} : & \quad \frac{ig}{2\sqrt{2} m_W} V_{tb} [m_b(1 + \gamma_5) \tan \beta + m_t(1 - \gamma_5) \cot \beta], \\ H^- t \bar{b} : & \quad \frac{ig}{2\sqrt{2} m_W} V_{tb}^* [m_b(1 - \gamma_5) \tan \beta + m_t(1 + \gamma_5) \cot \beta]. \end{aligned} \quad (2.13)$$

The $H_j ZZ$ ($H_j W^+ W^-$) coupling is, relative to that of the SM, given by

$$H_j ZZ \ (H_j W^+ W^-) : \quad [\cos \beta R_{j1} + \sin \beta R_{j2}], \quad \text{for } j = 1. \quad (2.14)$$

Note that when H_1 is CP-odd ($H_1 = A$), then $c_2 = 0$ and this vector coupling vanishes [26]. Finally, the $H_j H^+ W^-$ coupling is given by [35]

$$H_j H^\pm W^\mp : \quad \frac{g}{2} [\mp i(\sin \beta R_{j1} - \cos \beta R_{j2}) + R_{j3}] (p_\mu^j - p_\mu^\pm). \quad (2.15)$$

3 Constraining the parameter space

The multi-dimensional type-II 2HDM parameter space is severely restricted by a variety of theoretical and experimental constraints, which are discussed in the following.

3.1 Theoretical constraints

We impose three classes of theoretical constraints:

- **Positivity:** In order to have a stable potential, we impose positivity, $V(\Phi_1, \Phi_2) > 0$ as $|\Phi_1|, |\Phi_2| \rightarrow \infty$ [37–39]. Additionally, we must insist on a non-trivial solution to Eq. (2.11): $M_3^2 > 0$ and $M_2 \leq M_3$. While positivity may be satisfied for the given parameters of the potential, the considered minimum need not be the global one. However, it has been shown that, if a local charge-conserving minimum exists, then there can be no charge-breaking minimum [40–42]. Nevertheless, the potential of the 2HDM can have more than one charge-conserving minimum. We therefore check that the minimum obtained is the global one, following the approach of Ref. [43].
- **Tree-level unitarity:** We also impose tree-level unitarity on Higgs–Higgs scattering [44–48]. These conditions have a rather dramatic effect at “large” values of $\tan \beta$ and M_{H^\pm} , though some tuning of μ can extend the allowed range to larger values of $\tan \beta$ [49].

- **Perturbativity:** We impose the following upper bound on all λ 's:

$$|\lambda_i| < 4\pi\xi, \quad (3.1)$$

with $\xi = 0.8$, meaning $|\lambda_i| \lesssim 10$. The effect of this is to restrict large values of the masses, unless the soft parameter μ is comparable to M_2 and M_{H^\pm} .

For illustrations of how these theory constraints cut into the parameter space, see Refs. [34, 35].

3.2 Experimental constraints

Below, we list the different experimental constraints that are important. The SM predictions of the flavour observables quoted in this subsection are obtained using SuperIso v3.2 [50, 51].

- $B \rightarrow X_s \gamma$: This rare FCNC inclusive decay receives contributions from the charged Higgs boson that can be comparable to the W^\pm contribution in the SM. Since the charged Higgs state always contributes positively to the corresponding BR, it is an effective tool to probe the type-II 2HDM. The most up-to-date SM prediction for this decay, at the Next-to-Next-to-Leading Order (NNLO), gives [50–65]:

$$\text{BR}(\bar{B} \rightarrow X_s \gamma)_{\text{SM}} = (3.11 \pm 0.22) \times 10^{-4}, \quad (3.2)$$

while the combined experimental value from HFAG points to a larger value [66]:

$$\text{BR}(\bar{B} \rightarrow X_s \gamma)_{\text{exp}} = (3.55 \pm 0.24 \pm 0.09) \times 10^{-4}. \quad (3.3)$$

For type-II Yukawa interactions, which we consider here, light charged Higgs bosons are excluded by this observable. The actual limit is sensitive to higher-order QCD effects and is of the order of 300 GeV, being more severe at low values of $\tan \beta$ [59, 65, 67, 68]. Recently, a higher-order calculation [69] concludes that the 95% C.L. is at 380 GeV. While adopting the more conservative limit of 300 GeV, we shall also discuss this more restrictive one.

- $B_u \rightarrow \tau \nu_\tau$: In contrast to the $b \rightarrow s \gamma$ transitions, where the charged Higgs state participates in loop diagrams, the process $B_u \rightarrow \tau \nu_\tau$ can be mediated by H^\pm already at tree level. Since this decay is helicity suppressed in the SM, whereas there is no such suppression for spinless H^\pm exchange in the limit of large $\tan \beta$, these two contributions can be of similar magnitude [70]. The 2HDM contribution factorises in the ratio $R_{B\tau\nu}$ as compared to the SM value. This decay suffers from uncertainties from f_B and V_{ub} , and using $f_B = 192.8 \pm 9.9$ MeV [71], and the combined value $|V_{ub}| = (3.92 \pm 0.46) \times 10^{-3}$ [72], the SM BR evaluates numerically to [50, 51]:

$$\text{BR}(B_u \rightarrow \tau \nu_\tau)_{\text{SM}} = (1.01 \pm 0.29) \times 10^{-4}. \quad (3.4)$$

The SM prediction is compared to the current HFAG value¹ [66]

$$\text{BR}(B_u \rightarrow \tau \nu_\tau)_{\text{exp}} = (1.64 \pm 0.34) \times 10^{-4} \quad (3.5)$$

¹The latest BaBar result $\text{BR}(B_u \rightarrow \tau \nu_\tau) = (1.83^{+0.53}_{-0.49} \pm 0.24) \times 10^{-4}$ [73] is not included in this average.

by forming the ratio

$$R_{B\tau\nu}^{\text{exp}} \equiv \frac{\text{BR}(B_u \rightarrow \tau\nu_\tau)_{\text{exp}}}{\text{BR}(B_u \rightarrow \tau\nu_\tau)_{\text{SM}}} = 1.62 \pm 0.54. \quad (3.6)$$

In the framework of the 2HDM this leads to the exclusion of two sectors of the ratio $\tan\beta/M_{H^\pm}$ [65, 70, 74–78].

- **$B \rightarrow D\tau\nu_\tau$:** Compared to $B_u \rightarrow \tau\nu_\tau$, the semi-leptonic decays $B \rightarrow D\ell\nu$ have the advantage of depending on $|V_{cb}|$, which is known to greater precision than $|V_{ub}|$. In addition, the $\text{BR}(B \rightarrow D\tau\nu_\tau)$ is about 50 times larger than the $\text{BR}(B_u \rightarrow \tau\nu_\tau)$ in the SM. The experimental determination remains however very complicated due to the presence of at least two neutrinos in the final state. The ratio

$$\xi_{D\ell\nu_\tau} = \frac{\text{BR}(B \rightarrow D\tau\nu_\tau)}{\text{BR}(B \rightarrow D\ell\nu_\ell)}, \quad (3.7)$$

where the 2HDM contributes only to the numerator, allows one to reduce some of the theoretical uncertainties. The SM prediction for this ratio is [50, 51]

$$\xi_{D\ell\nu}^{\text{SM}} = (29.7 \pm 3) \times 10^{-2}, \quad (3.8)$$

and the most recent experimental result by the BaBar collaboration is [79]

$$\xi_{D\ell\nu}^{\text{exp}} = (44.0 \pm 5.8 \pm 4.2) \times 10^{-2}. \quad (3.9)$$

This ratio is also sensitive to a light charged Higgs boson, and leads to complementary constraints to the ones following from $B_u \rightarrow \tau\nu_\tau$ [65, 78, 80–82].

- **$D_s \rightarrow \tau\nu_\tau$:** Constraints on a light charged Higgs can be obtained, competitive with those obtained from $B_u \rightarrow \tau\nu_\tau$ [83]. The main uncertainty here is due to the decay constant f_{D_s} . The SM prediction for this decay is [50, 51]:

$$\text{BR}(D_s \rightarrow \tau\nu_\tau)_{\text{SM}} = (5.11 \pm 0.13) \times 10^{-2}, \quad (3.10)$$

using $f_{D_s} = 248 \pm 2.5$ MeV [84], and the current world average of the experimental measurements gives [66]:

$$\text{BR}(D_s \rightarrow \tau\nu_\tau)_{\text{exp}} = (5.38 \pm 0.32) \times 10^{-2}. \quad (3.11)$$

- **$B_{d,s} \rightarrow \mu^+\mu^-$:** These decays are helicity suppressed in the SM and can receive sizeable enhancement or depletion from Higgs-mediated contributions. At large $\tan\beta$, the non-observation of these decay modes imposes a lower bound on the charged Higgs boson mass [85, 86]. The most stringent limits for their BRs were reported very recently by the LHCb collaboration [87]:

$$\text{BR}(B_s \rightarrow \mu^+\mu^-) < 4.5 \times 10^{-9}, \quad (3.12)$$

$$\text{BR}(B_d \rightarrow \mu^+\mu^-) < 1.0 \times 10^{-9}, \quad (3.13)$$

at 95% C.L. Combining these values with the limits from ATLAS and CMS [88] results in even stronger bounds:

$$\text{BR}(B_s \rightarrow \mu^+ \mu^-) < 4.2 \times 10^{-9}, \quad (3.14)$$

$$\text{BR}(B_d \rightarrow \mu^+ \mu^-) < 8.1 \times 10^{-10}. \quad (3.15)$$

The SM predictions for these branching ratios are [50, 51, 89]:

$$\text{BR}(B_s \rightarrow \mu^+ \mu^-) = (3.53 \pm 0.38) \times 10^{-9}, \quad (3.16)$$

$$\text{BR}(B_d \rightarrow \mu^+ \mu^-) = (1.1 \pm 0.1) \times 10^{-10}, \quad (3.17)$$

with $\text{BR}(B_s \rightarrow \mu^+ \mu^-)$ being the more constraining. The largest uncertainty is from f_{B_s} , we used $f_{B_s} = 234 \pm 10$ MeV for our evaluation. In the type-II 2HDM, the experimental limits can be reached for very large values of the Yukawa couplings and small charged Higgs boson masses. The constraining power of $B_{d,s} \rightarrow \mu^+ \mu^-$ in this study is hence rather limited as compared to the other flavour observables.

- **$B^0 - \bar{B}^0$ mixing:** Due to the possibility of H^\pm exchange, in addition to W exchange, the $B^0 - \bar{B}^0$ mixing constraint, which is sensitive to the term $m_t \cot \beta$ in the Yukawa couplings (2.13), excludes low values of $\tan \beta$ and low values of M_{H^\pm} [65, 90–95]. The non-perturbative decay constant f_{B_d} and the bag parameter \hat{B}_d which are evaluated simultaneously from lattice QCD constitute the largest theoretical uncertainty.
- **R_b :** The branching ratio $R_b \equiv \Gamma_{Z \rightarrow b\bar{b}}/\Gamma_{Z \rightarrow \text{had}}$ would also be affected by Higgs boson exchange. The contributions from neutral Higgs bosons to R_b are negligible [35], however, charged Higgs boson contributions, via the $H^\pm b t$ Yukawa coupling, as given by [96], Eq. (4.2), exclude low values of $\tan \beta$ and low M_{H^\pm} .
- **$pp \rightarrow H_j X$:** Two aspects of the recent neutral Higgs searches at the LHC are considered [97, 98]:
 - The production and subsequent decay of a neutral Higgs to $\gamma\gamma$, around $M = 125$ GeV is taken to be within a factor of 2 from the SM rate. Assuming the dominant production to be via gluon fusion, this can be approximated as $0.5 \leq R_{\gamma\gamma} \leq 2$, where we define

$$R_{\gamma\gamma} = \frac{\Gamma(H_1 \rightarrow gg)\text{BR}(H_1 \rightarrow \gamma\gamma)}{\Gamma(H_{\text{SM}} \rightarrow gg)\text{BR}(H_{\text{SM}} \rightarrow \gamma\gamma)}. \quad (3.18)$$

We take into account (1) the modified scalar coupling to the fermion or W in the loop, (2) the pseudoscalar contribution, and (3) the charged Higgs contribution on the $\gamma\gamma$ side. This condition (3.18) mainly constrains the Yukawa couplings of H_1 . In particular, the (dominant) $H_1 t\bar{t}$ contribution to the loop integrals should be comparable to that of the SM, meaning

$$\frac{s_1^2 c_2^2}{\sin^2 \beta} + \frac{s_2^2}{\tan^2 \beta} P^2(\tau_t) = \mathcal{O}(1), \quad (3.19)$$

where $P(\tau)$ represents the ratio of the pseudoscalar and the scalar contributions to the loop integral [6], with $\tau_t = M_1^2/(4m_t^2)$. At low $\tan\beta$ there is some freedom, either $s_1^2 c_2^2$ or s_2^2 should be of order unity, whereas at high $\tan\beta$ this constraint requires $\alpha_1 \simeq \pm\pi/2$ and $\alpha_2 \simeq 0$.

- The production and subsequent decay, dominantly via ZZ and WW , is constrained in the mass range $130 \text{ GeV} \lesssim M \lesssim 600 \text{ GeV}$. We consider the quantity

$$R_{ZZ} = \frac{\Gamma(H_j \rightarrow gg)\text{BR}(H_j \rightarrow ZZ)}{\Gamma(H_{\text{SM}} \rightarrow gg)\text{BR}(H_{\text{SM}} \rightarrow ZZ)}, \quad (3.20)$$

for $j = 2, 3$ and require it to be below the stronger 95% CL obtained by ATLAS or CMS. This constraint thus affects the product of the Yukawa and gauge couplings of H_2 and H_3 (see Eqs. (2.12) and (2.14)). In the limit of unchanged total width, this implies

$$\left[\frac{R_{j2}^2}{\sin^2 \beta} + \frac{R_{j3}^2}{\tan^2 \beta} \right] [\cos \beta R_{j1} + \sin \beta R_{j2}]^2 < \eta, \quad j = 2, 3, \quad (3.21)$$

where η is the 95% CL on $\sigma/\sigma_{\text{SM}}$. For $\tan\beta$ of the order of unity, the first factor is “small” when R_{j1}^2 is of order unity, whereas the second factor is “small” when $|R_{j3}|$ is of order unity.

For larger values of $\tan\beta$, we may substitute $\alpha_1 \simeq \pm\pi/2$ and $\alpha_2 \simeq 0$ from the above consideration, whereupon the constraint (3.21) takes the form:

$$\frac{\cos^2 \beta}{\tan^2 \beta} (s_3 c_3)^2 < \eta, \quad (3.22)$$

which is not very strong. In particular, it is automatically satisfied at large $\tan\beta$.

For the total widths of H_2 and H_3 entering in the numerator of (3.20), we include also $H_j \rightarrow H_1 H_1$ and $H_j \rightarrow H_1 Z$, in addition to the familiar decay modes of H_1 . The relevant couplings can be found in [99] and [35].

- **T and S :** For the electroweak “precision observables” T and S , we impose the bounds $|\Delta T| < 0.10$, $|\Delta S| < 0.10$ [100], at the 1- σ level, within the framework of Refs. [101, 102]. While S is not very restrictive, T gets a positive contribution from a splitting between the masses of charged and neutral Higgs bosons, whereas a pair of neutral ones give a negative contribution.
- **Electron Electric Dipole Moment (EDM):** The bound on the electron EDM constrains the allowed amount of CP violation of the model. We adopt the bound [103] (see also [104]):

$$|d_e| \lesssim 1 \times 10^{-27} [e \text{ cm}], \quad (3.23)$$

at the 1- σ level. The contribution due to neutral Higgs exchange, via the two-loop Barr–Zee effect [105], is given by Eq. (3.2) of [104].

In contrast to the MSSM, in the 2HDM, an additional contribution to the muon anomalous magnetic moment arises only at the two-loop level. Since we are considering heavy Higgs bosons ($M_1, M_{H^\pm} \gtrsim 100$ GeV) therefore, according to [68, 106, 107], the 2HDM contribution to the muon anomalous magnetic moment is negligible even for $\tan \beta$ as large as ~ 40 .

The above constraints are not independent. Therefore, we do not attempt to add their contributions to an overall χ^2 , but rather require that none of them should be violated by more than 2σ . The LHC constraints are imposed at the quoted 95% C.L.

Among these constraints, basically $b \rightarrow s\gamma$ requires $M_{H^\pm} \gtrsim 300$ (380) GeV, while the different B -meson constraints impose additional constraints at low and high values of $\tan \beta$ (the latter are basically excluded anyway, by the unitarity constraints). The T constraint prevents the masses of the neutral and charged Higgs bosons from being very different, and thus effectively provides a cut-off at high masses.

3.3 Two scenarios

In order to develop some intuition for the viable parts of the parameter space, we shall here consider two scenarios. In both of them, we take the lightest neutral Higgs boson mass to be $M_1 = 125$ GeV. Furthermore:

- **Scenario 1: low $\tan \beta$, intermediate different masses M_2 and M_{H^\pm} (non-decoupling regime):**

$$\begin{aligned} \tan \beta = 1, 2, 3; \quad M_1 = 125 \text{ GeV}, \quad M_2 = 150, 200, 300, 400, 500 \text{ GeV}, \\ M_{H^\pm} = 300 - 600 \text{ GeV}, \quad \mu = 200 \text{ GeV}. \end{aligned} \quad (3.24)$$

- **Scenario 2: high $\tan \beta$, heavy degenerate masses M_2 and M_{H^\pm} (decoupling regime):**

$$\begin{aligned} \tan \beta = 5, 10, 20; \quad M_1 = 125 \text{ GeV}, \quad M_2 = 400, 600 \text{ GeV}, \\ M_{H^\pm} \simeq M_2, \quad \mu \simeq M_2. \end{aligned} \quad (3.25)$$

For scenario 1, we will consider a range of charged Higgs boson masses, from 300 to 600 GeV (for 700 GeV, only a few viable parameter points are found), and a range of $\tan \beta$ values. We will typically consider small values of $\tan \beta$, of the order of 1. High values lead (for fixed μ) to conflict with the unitarity constraints.

Scenario 2 is rather fine-tuned. The masses M_2 and M_{H^\pm} have to be very close to μ , in order to avoid conflict with the unitarity constraints. It is discussed in some detail in the appendix. There, it is shown that in addition to the SM-like region of $\alpha_1 \simeq \pm\pi/2$ with $\alpha_2 \simeq 0$ (H_2 or H_3 being odd under CP), there is also another region, with $\alpha_1 \simeq 0$ and $\alpha_2 \simeq \pm\pi/2$ (with H_1 being odd, see [108, 109]).

3.4 Studies in α space

3.4.1 Scenario 1

We shall first consider separately the LHC constraints on $H_1 \rightarrow \gamma\gamma$ (an allowed range) and $H_{2,3} \rightarrow W^+W^-$ (an upper bound). For fixed additional input parameters,

$$M_{H^\pm} = 300, 500 \text{ GeV}, \quad \tan\beta = 1, \quad \mu = 200 \text{ GeV}, \quad (3.26)$$

we show in fig. 1 the result of a scan over 10^6 sets of $(\alpha_1, \alpha_2, \alpha_3)$. The blue regions satisfy all the constraints discussed above, as well as one of these LHC constraints, whereas the red regions do not satisfy the LHC constraint considered, but all other constraints discussed in sections 3.1 and 3.2.

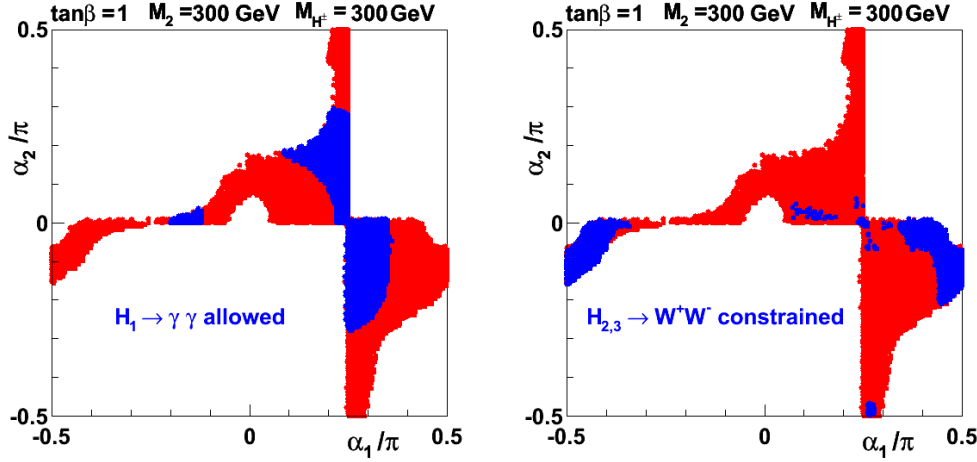


Figure 1: Allowed regions in the α_1 - α_2 parameter space, without (red) and with (blue) LHC constraints, for $M_2 = 300$ GeV and the additional parameters given in Eq. (3.26). Left, blue: regions surviving the $H_1 \rightarrow \gamma\gamma$ constraint. Right, blue: regions surviving the $H_{2,3} \rightarrow W^+W^-$ constraint.

The allowed regions are for these parameters rather independent of M_{H^\pm} in the range 300–400 GeV, but start shrinking around 500 GeV and vanish around $M_{H^\pm} \sim 600$ GeV. (If we allow $\xi = 1$, see Eq. (3.1), the allowed values of M_{H^\pm} reach out to about 700 GeV.) The underlying checkered pattern is due to the positivity constraint, together with $M_3^2 \geq M_2^2$.

When we impose the LHC constraints discussed above [97, 98] (scanning now over $5 \cdot 10^6$ points), we obtain the blue regions for the effects of the $H_1 \rightarrow \gamma\gamma$ (left) and the $H_{2,3} \rightarrow W^+W^-$ (right) constraint. For the considered case of $\tan\beta = 1$, $M_2 = M_{H^\pm} = 300$ GeV we see that these only very marginally overlap.

Imposing then both of these LHC constraints, we show in figure 2 the resulting surviving parameter space in green on top of the red regions, now for $M_2 = 400$ GeV, and two values of charged-Higgs mass, 300 and 500 GeV. Note that although the LHC experiments exclude an SM Higgs with a mass from about 130 GeV to about 600 GeV, there are still viable regions of parameter space for the second (and third) Higgs state to be in this region, since it may couple more weakly than the SM Higgs boson.

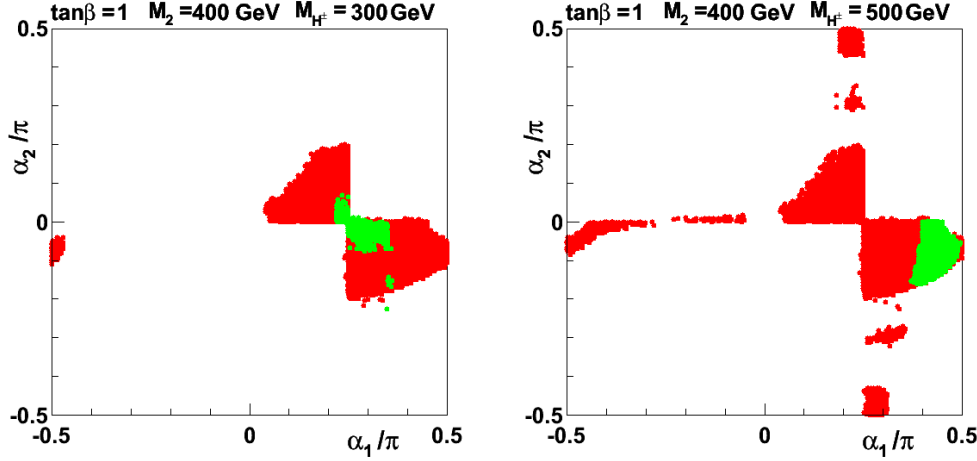


Figure 2: Allowed regions in the α_1 - α_2 parameter space, without (red) and with (green) the LHC constraints, for $M_2 = 400$ GeV and the additional parameters given in Eq. (3.26).

A striking first observation is that the allowed regions are very much reduced, only values close to $\alpha_1 = \pm\pi/4$ are now allowed for the low value of M_{H^\pm} , and a somewhat higher value for the higher value of M_{H^\pm} . This is a little different from the results reported recently for the CP-conserving model, where it was found that only a region around $\alpha = 0$ is allowed [24]. We recall that in the particular CP-conserving limit of $\alpha_2 = \alpha_3 = 0$ (corresponding to the heaviest one, H_3 , being odd under CP), $\alpha = 0$ corresponds to $\alpha_1 = \pm\pi/2$. That region is here found to violate unitarity.

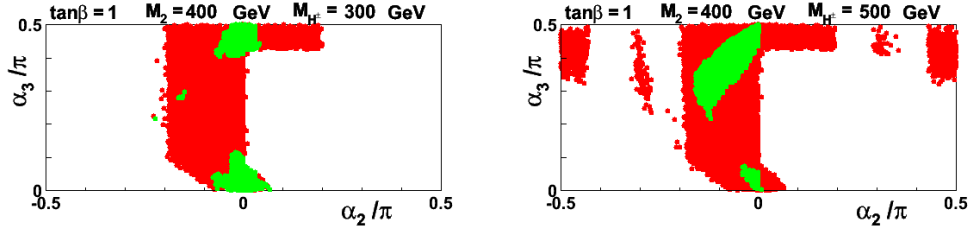


Figure 3: Allowed regions in the α_2 - α_3 parameter space, without (red) and with (green) the LHC constraints, for the parameters given in Eq. (3.26).

For a complementary view of the allowed region, we show in fig. 3 the corresponding projections onto the α_2 - α_3 plane. Much of the region with $\alpha_2 < 0$ is populated, whereas most of the region with $\alpha_2 > 0$ is excluded. All values of α_3 are represented. When we impose the LHC constraints, the main characteristic is that α_2 becomes more restricted, as was also seen in fig. 2, and by Barroso et al [27]. Two regions of α_3 -values are represented: values close to 0 or $\pi/2$ (in both cases for α_2 close to zero).

A comparison with Arhrib et al [28] indicates that we find a more constrained region, presumably because of our tighter constraint on $R_{\gamma\gamma}$. In addition, we find that some points

are excluded because of conflict with the electron EDM.

For higher values of $\tan \beta$, several things change. At some point, also negative values of α_1 become allowed (not shown), and the allowed ranges of α_1 move towards $\pm\pi/2$ as $\tan \beta$ increases. At $\tan \beta = 2$ and 3, the allowed region has shrunk by a factor of 2 to 3, compared to $\tan \beta = 1$.

For lower values of M_2 , the impact of the LHC constraints is more severe, but allowed points are found, for example also for $M_2 = 150$ GeV. For higher values of M_2 , the allowed region is restricted to some neighborhood of $\alpha_1 = \pi/4$ and $\alpha_2 = 0$, with α_3 close to 0 or $\pi/2$.

In view of the results shown in fig. 2, let us comment on the special limit

$$\alpha_1 = \pi/4, \quad \alpha_2 \text{ small } (s_2 \simeq 0). \quad (3.27)$$

Then, the rotation matrix can be simplified as

$$R = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 1 & 0 \\ -c_3 & c_3 & \sqrt{2}s_3 \\ s_3 & -s_3 & \sqrt{2}c_3 \end{pmatrix}, \quad (3.28)$$

and the physical states are related to the “weak” states η_j as

$$H_1 = \frac{1}{\sqrt{2}}(\eta_1 + \eta_2), \quad (3.29a)$$

$$H_2 = \frac{1}{\sqrt{2}}(-c_3\eta_1 + c_3\eta_2 + \sqrt{2}s_3\eta_3), \quad (3.29b)$$

$$H_3 = \frac{1}{\sqrt{2}}(s_3\eta_1 - s_3\eta_2 + \sqrt{2}c_3\eta_3). \quad (3.29c)$$

We recall that, with type-II Yukawa couplings, η_1 couples to down-type quarks whereas η_2 couples to up-type quarks. Thus, in the limit (3.27), the lightest neutral Higgs boson couples coherently to the b - and the t -quarks (with strengths proportional to $1/\cos \beta$ and $1/\sin \beta$, respectively, and thus, for $\tan \beta = 1$, with the same strength). The heavier Higgs bosons, however, have a CP-even content given by $\eta_1 - \eta_2$ (note the minus sign). For $\tan \beta = 1$ only the pseudoscalar component, η_3 , of these heavier Higgs bosons will couple to fermions. We note that two CP-conserving limits are contained in this scenario: $H_3 = A$ for $\alpha_3 = 0$ and $H_2 = A$ for $\alpha_3 = \pi/2$.

3.4.2 Scenario 2

Again, we start out with an overview of the allowed regions of parameter space in the absence of LHC constraints. This is presented in red, in fig. 4, for $\tan \beta = 10$ and two values of M_2 , namely 400 (left) and 600 (right) GeV. The populated regions are at $(\alpha_1, \alpha_2) \simeq (\pm\pi/2, 0)$ and $\simeq (0, \pm\pi/2)$, the decoupling regions which are discussed in the appendix. They are seen to shrink considerably as the masses are increased from 400 to 600 GeV.

When we impose the LHC constraints, the regions near $(\alpha_1, \alpha_2) = (0, \pm\pi/2)$ are no longer allowed. Neither is the region near $(\alpha_1, \alpha_2) \simeq (-\pi/2, 0)$. This is a CP-conserving limit, also commented on in section 3.4.3. Only a small region near $(\alpha_1, \alpha_2) \simeq (\pi/2, 0)$

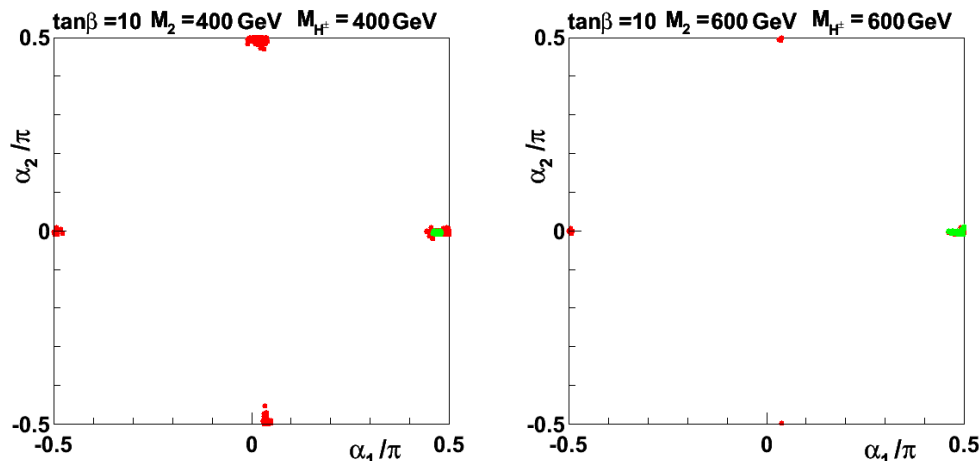


Figure 4: Allowed regions in the α_1 - α_2 parameter space, without (red) and with (green) the LHC constraints, for $M_1 = 125$ GeV, $M_2 = 400$ GeV (left), $M_2 = 600$ GeV (right) and $M_2 - 10$ GeV $\leq M_{H^\pm} \leq M_2 + 10$ GeV.

remains. The figure includes a range of values for M_{H^\pm} within $[M_2 - 10$ GeV, $M_2 + 10$ GeV]. If we take $M_{H^\pm} = M_2$, then α_2 has to be slightly different from zero, meaning that CP is violated. At the CP-conserving parameter point $(\alpha_1, \alpha_2, \alpha_3) = (\pi/2, 0, 0)$, unitarity is violated.

3.4.3 CP-conserving limits

Solutions also exist in CP-conserving limits. With H_3 CP-odd ($H_3 = A$), we have performed scans at $\tan \beta = 3, 5, 10$ and 20 , for two heavy-mass cases: (i) $H_2 = M_{H^\pm} \equiv M = 400$ GeV and (ii) $M = 600$ GeV. In the absence of the LHC constraint, some range in α_1 around $\pm\pi/2$ (corresponding to α around 0) is populated. Imposing the LHC constraints, the range in α_1 is constrained, see fig. 5 for the case $\tan \beta = 5$. This is consistent with the allowed regions discussed in section 3.4.1, for $\alpha_2 \simeq \alpha_3 \simeq 0$ (see figures 2 and 3). For increasing values of $\tan \beta$ (10 and 20), the allowed bands move towards the edges and become more narrow (compare fig. 4).

When H_2 is CP-odd ($H_2 = A$) and $\tan \beta = 5$, we find allowed solutions at $M = 400$ GeV and 600 GeV. But compared to the case $H_3 = A$, the mass range for M_3 is more constrained.

We do not find any solution when H_1 is CP-odd ($H_1 = A$). The crucial LHC constraint is the quantity $R_{\gamma\gamma}$. Others [25, 26] have argued for an interpretation of the 125 GeV excess in terms of a pseudoscalar, but we are not able to confirm this. We find points which satisfy all other constraints, but not the LHC constraints. Another recent study agrees with this [110].

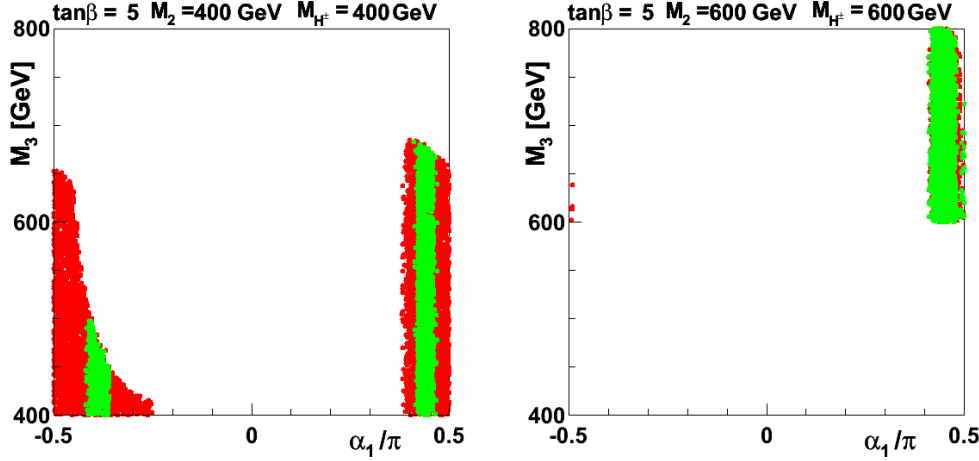


Figure 5: CP conserving case. Allowed regions in the α_1 – M_3 parameter space, without (red) and with (green) the LHC constraints, for $M_1 = 125$ GeV, $M_2 = M_{H^\pm} = 400$ GeV (left), $M_2 = M_{H^\pm} = 600$ GeV (right).

3.5 Studies in $\tan\beta$ – M_{H^\pm} space

3.5.1 The unitarity constraint

The unitarity constraint plays an important role in delimiting high values of both $\tan\beta$ and M_{H^\pm} . This constraint requires the λ 's to be small, which to some extent is achieved by taking the “soft” mass parameter μ large. In fact, in the co-called decoupling limit, discussed for the CP-conserving case in [111], and for the present case in the appendix, one can respect the unitarity constraints for large masses, provided μ is tuned to these masses:

$$M_2 \sim M_3 \sim M_{H^\pm} \sim \mu. \quad (3.30)$$

For moderate values of $\tan\beta$ ($3 - 5$), that limit also requires

$$\beta \sim \alpha_1, \quad \alpha_2 \sim 0, \quad \alpha_3 \text{ arbitrary}. \quad (3.31)$$

For large values of $\tan\beta$ ($\gtrsim 5$), this evolves into the region $(\alpha_1, \alpha_2) \sim (\pi/2, 0)$. Furthermore, an additional region opens up for large masses and large $\tan\beta$, leading to

$$\text{Decoupling 1: } (\alpha_1, \alpha_2) \sim (\pm\pi/2, 0), \quad (3.32a)$$

$$\text{Decoupling 2: } (\alpha_1, \alpha_2) \sim (0, \pm\pi/2), \quad (3.32b)$$

with α_3 arbitrary. Two comments are here in order: (i) because of the periodicity of the trigonometric functions, regions at $\alpha_i \simeq -\pi/2$ and $\alpha_i \simeq +\pi/2$ are connected; (ii) the SM limit requires $\alpha_1 \sim \beta$, and is thus contained in the region “Decoupling 1”.

In view of the above discussion, in order to determine the maximally allowed ranges of $\tan\beta$ and M_{H^\pm} , we scan over some range in μ , starting at the geometric mean

$$\mu_0 = \sqrt{M_2 M_{H^\pm}}. \quad (3.33)$$

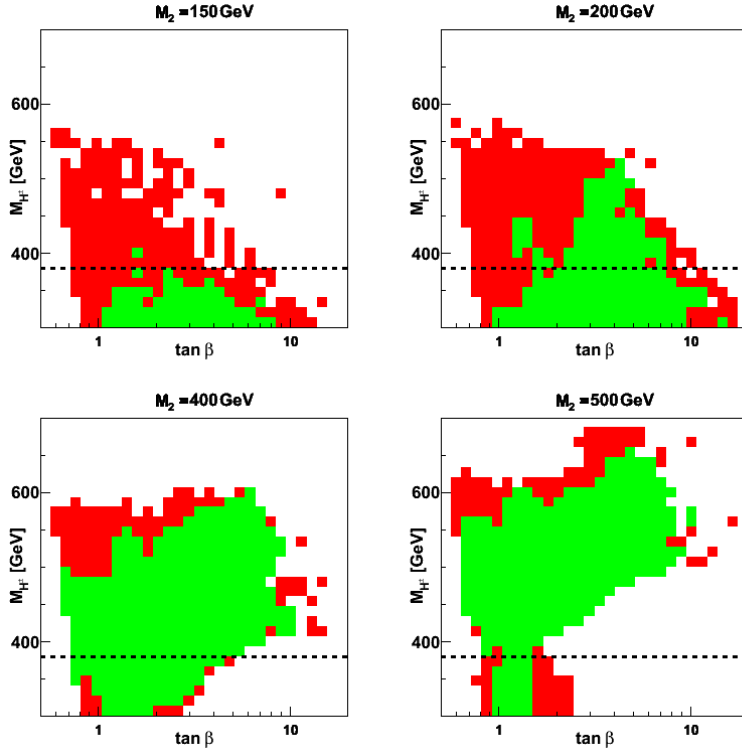


Figure 6: Allowed regions in the $\tan\beta$ – M_{H^\pm} parameter space, without (red) and with (green) the LHC constraints, for $M_1 = 125$ GeV and four values of M_2 , as indicated. The dashed lines show the recent bound at 380 GeV [69].

3.5.2 The experimental constraints

In fig. 6 we show allowed regions in the $\tan\beta$ – M_{H^\pm} plane. Again, the larger red region is allowed in the absence of recent LHC results, whereas the green region shows what remains compatible with these data. We note some reduction in the range of charged Higgs masses. Also, at high $\tan\beta$, the masses M_2 and M_{H^\pm} tend to be close, as discussed above.

The “fractal” appearance of these plots is in part due to the finite number of points in the scans. Some could also reflect genuine “islands” in parameter space.

In fig. 7 we show typical values of M_3 . Note that for each point in the allowed part of this plane, some ranges of α ’s are allowed (see the previous subsection). Each set of α ’s corresponds to a particular value of M_3 . The values plotted here are those first encountered in a random scan over α ’s. We see that as M_2 and M_{H^\pm} increase, also typical values of M_3 increase.

By allowing a larger value of the perturbativity cut-off ξ of Eq. (3.1), higher masses of M_{H^\pm} would be allowed. For example, $\xi = 1$ permits masses above 600 GeV. Also the unitarity and the electroweak parameter T constrain this high-mass region. Which of these gives the strongest limits depends on the other parameters.

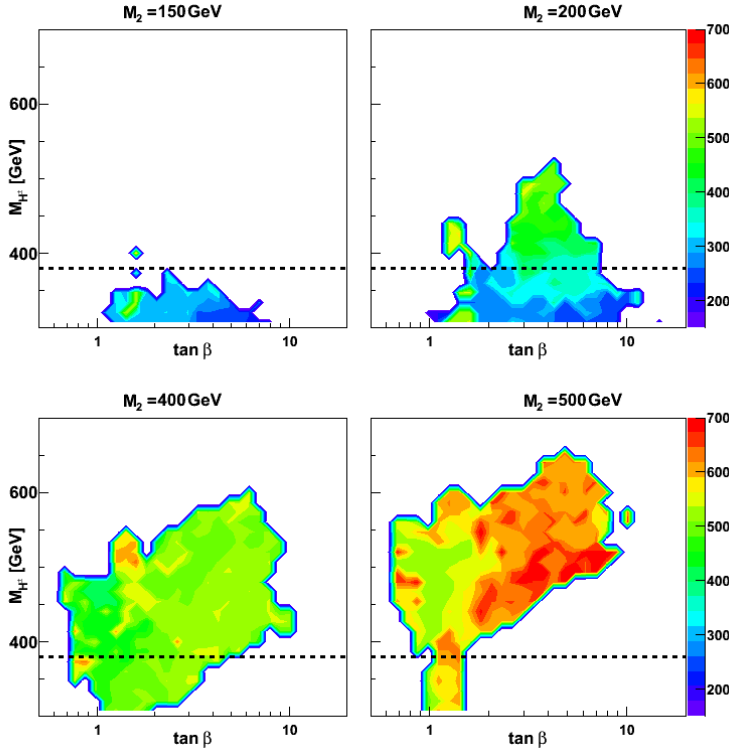


Figure 7: Heaviest neutral Higgs mass, M_3 in the $\tan\beta$ – M_{H^\pm} parameter space, with the LHC constraints, for $M_1 = 125$ GeV and four values of M_2 , as indicated. The dashed lines show the recent bound at 380 GeV [69].

4 Benchmark analysis

In this section we study the profile of the charged Higgs boson at the LHC in view of the allowed parameter space analysis. For this, we start by studying a set of candidate benchmark points P_i that allow us to synthesise the main features of the surviving models.

As we have shown in the previous sections, this model depends on eight parameters. However, since the μ parameter does not directly participate in the phenomenology of interest, if not specified otherwise, we will consider $\mu = 200$ GeV hereafter. The exception will be the high- $\tan\beta$ case, where μ has to be carefully tuned in order to respect the unitarity constraint. Then, we consider points that pass the constraints, at least in the charged Higgs mass range 300 – 600 GeV. We are then left with five parameters: $\tan\beta$, α_1 , α_2 , α_3 and M_2 .

In table 1, we list a set of 10 candidate benchmark points: we consider the most illustrative four of them for determining cross sections and relevant BRs, the rest will be discussed only qualitatively.

This set of points has distinctive characteristics in the phenomenology, as will be shortly made clear. Also, we note that P_9 and P_{10} are nearly CP-conserving, with $\alpha_2 \simeq \alpha_3 \simeq 0$. All the others correspond to CP-violating scenarios.

	α_1/π	α_2/π	α_3/π	$\tan \beta$	M_2	$M_{H^\pm}^{\min}, M_{H^\pm}^{\max}$
P_1	0.23	0.06	0.005	1	300	300,325
P_2	0.35	-0.014	0.48	1	300	300,415
P_3	0.35	-0.015	0.496	1	350	300,450
P_4	0.35	-0.056	0.43	1	400	300,455
P_5	0.33	-0.21	0.23	1	450	300,470
P_6	0.27	-0.26	0.25	1	500	300,340
P_7	0.39	-0.07	0.33	2	300	300,405
P_8	0.34	-0.03	0.11	2	400	300,315
P_9	0.47	-0.006	0.05	10	400	400,440
P_{10}	0.49	-0.002	0.06	10	600	600,700

Table 1: Benchmark points selected from the allowed parameter space. Masses M_2 and allowed range of M_{H^\pm} are in GeV. For P_1 – P_8 , $\mu = 200$ GeV, whereas for P_9 and P_{10} , $\mu = M_2$.

The model has been implemented through the LanHEP module [112] (see [113] for details) and the following analysis has been performed by means of the CalcHEP package [114, 115]. Furthermore, we have used the CTEQ6.6M [116] set of five-flavour parton distribution functions (PDFs). Due to their relevance at hadron colliders, the effective ggH_i , $\gamma\gamma H_i$ and $\gamma Z H_i$ vertices have been implemented by means of a link between CalcHEP and LoopTools [117], and the numerical results have been cross-checked against the analytical results in [118].

4.1 Charged and lightest neutral Higgs bosons: BRs

This subsection is devoted to an analysis of the BRs of the charged Higgs state in the allowed parts of the parameter space. In fact, since we are mainly interested in signatures of a charged Higgs boson produced in association with a W boson, which involve model dependent couplings, it is of fundamental importance to establish some characteristic features of the BRs for some specific points of parameter space. In this regard, we consider four points from table 1 and we determine the most important decay modes. We consider only BRs $> 10^{-4}$, rates below this value are not of phenomenological relevance. Then, we have six decay modes: WH_1 , WH_2 , WH_3 , tb , ts , $\tau\nu_\tau$, displayed in fig. 8 for selected benchmark points. In addition, we remark that the results are presented for an illustrative range of $M_H = 300 - 600$ GeV, while recalling that the allowed range is always bounded in the range $M_H \sim 380 - 450$ GeV for an intermediate choice of $\tan \beta$.

First, we consider two points associated to the choice of $\tan \beta = 1$ in fig. 8, P_2 and P_6 . The dominant decay mode is always tb , and this feature is even reinforced when the masses M_1 and M_2 are closer (P_2), with respect to the case in which they are well separated (P_6). However, it is important to remark that the WH_i branching fractions, when allowed by phase space, are always $\sim \mathcal{O}(0.1)$ and never smaller than $\sim \mathcal{O}(0.01)$. In particular, if $M_{H^\pm} > 400$ GeV, then the BR for WH_1 is $\sim \mathcal{O}(0.1)$, this assures that the suppression

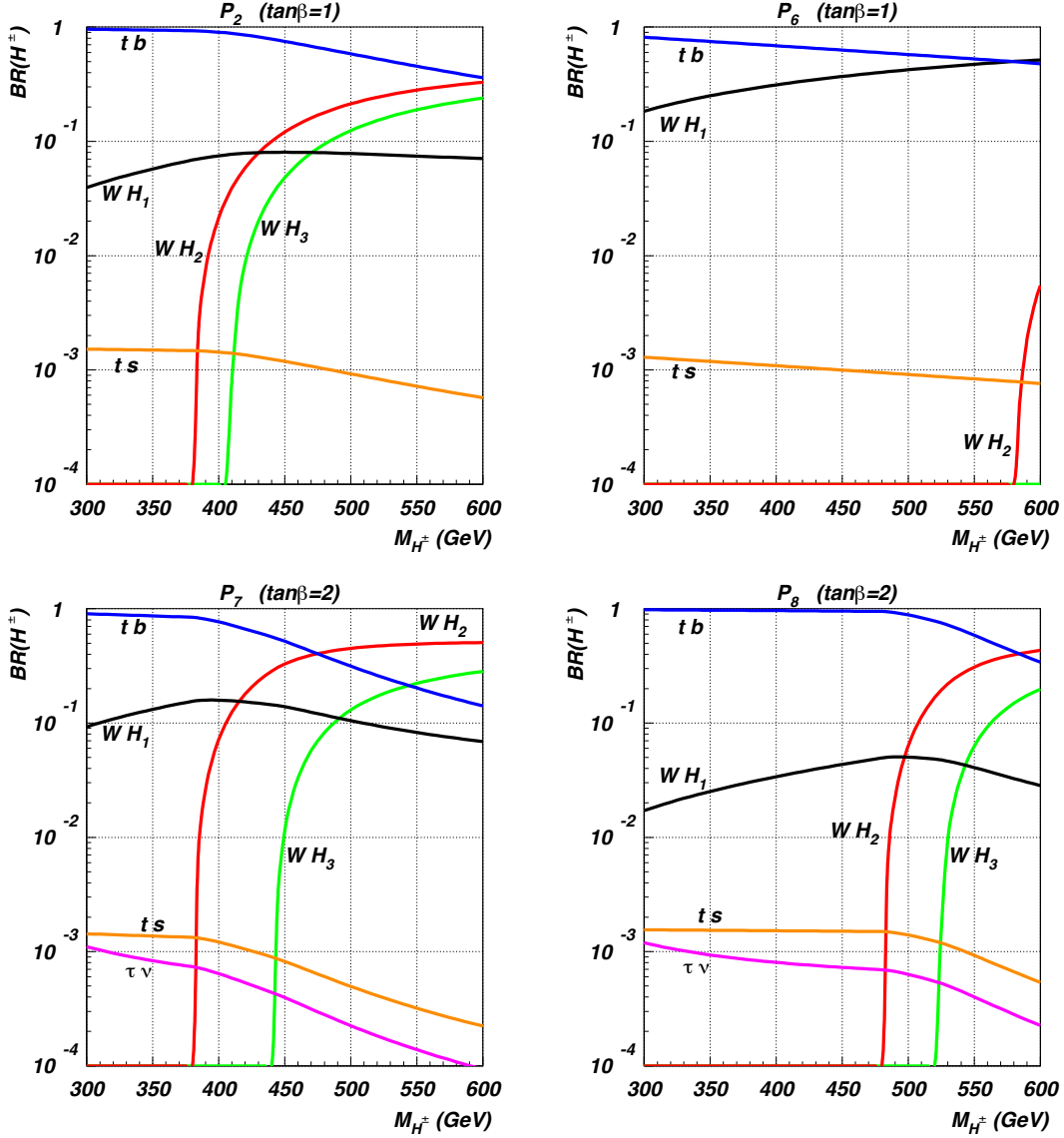


Figure 8: Branching ratios of the charged Higgs vs mass. Four benchmark points have been considered, as indicated.

brought about by this decay mode is never stronger than about an order of magnitude for a rather large value of M_{H^\pm} .

The result does not hold for the $W^\pm H_1$ case when $\tan\beta = 2$ (see fig. 8, P_7 and P_8). In fact, it strongly depends on the choice of point in parameter space: for P_8 this decay mode is suppressed down to $\sim \mathcal{O}(0.01)$, while for P_7 its branching fraction is restored to $\sim \mathcal{O}(0.1)$ because increasing the mixing between η_1 and η_3 via $|\alpha_2/\pi| = -0.03 \rightarrow -0.07$ increases its CP-odd component, and this effect leads to an enhancement.

Another feature of the $\tan\beta = 2$ choice is that the $W^\pm H_2$ decay mode is always dominant as compared to the tb one when M_{H^\pm} is large (though not always allowed),

H_1 decay modes	P_2	P_6	P_7	P_8
$b\bar{b}$	0.3414	0.5916	0.2595	0.3697
$s\bar{s}$	0.0001	0.0002	0.0000	0.0001
$c\bar{c}$	0.0625	0.0317	0.0805	0.0575
$\tau^+\tau^-$	0.0416	0.0721	0.0316	0.0451
$\mu^+\mu^-$	0.0001	0.0002	0.0001	0.0002
W^+W^-	0.3158	0.1241	0.3218	0.3051
gg	0.1944	0.1621	0.2617	0.1796
ZZ	0.0386	0.0152	0.0393	0.0373
γZ	0.0024	0.0010	0.0024	0.0023
$\gamma\gamma$	0.0032	0.0018	0.0030	0.0031

Table 2: BRs of the lightest neutral Higgs, $\text{BR}(H_1 \rightarrow X)$, for a set of benchmark points extracted from table 1: $P_1, P_5, P_7, P_8, P_9, P_{10}$.

due to the suppression of the $H^\pm \rightarrow tb$ coupling by a factor ~ 2 plus an always sizeable $H^\pm \rightarrow W^\pm H_2$ coupling. We remark that in this scenario the WH_1 BR is anyway $\sim \mathcal{O}(0.01)$ or bigger².

Since we are interested in the phenomenology of the charged Higgs boson produced in association to vector bosons, it is important to understand the properties of the lightest neutral Higgs. For this, we conclude this subsection by presenting a relevant set of H_1 BRs. As is clear from table 2, the three most important decay mode are always the $b\bar{b}$, gg and WW^* ones. Since the first is the phenomenologically simplest among the three, we will only consider this decay channel for studying combined H_1 signatures.

4.2 Charged Higgs: single production mechanisms

In this subsection we study single charged Higgs production at the LHC for two choices of energy, $\sqrt{s} = 8$ TeV and $\sqrt{s} = 14$ TeV. Considering the partonic amplitudes, we have three main mechanisms to produce a single charged Higgs boson from hadrons, i.e., associated with bosons (**B**) or fermions (**F**):

$$\mathbf{B}: \quad (gg, b\bar{b}) \rightarrow H^\pm W^\mp \quad (\text{fig. 9a}); \quad (4.1)$$

$$\mathbf{B}: \quad q\bar{q}' \rightarrow H^\pm H_i \quad (\text{fig. 9b}); \quad (4.2)$$

$$\mathbf{F}: \quad gg \rightarrow H^\pm b\bar{t} \quad (\text{fig. 9c}). \quad (4.3)$$

In fig. 9 we show the main partonic contributions to the three production channels.

Note that the process in fig. 9b will generally be disfavoured for two reasons. First, the quark density inside the proton is lower than the gluon density, so this channel is suppressed in this respect, owing to the typical H^\pm masses considered (recall that $x^2 \propto M_{H^\pm}^2/s$, where

²The $\tau\nu_\tau$ decay mode, on the other hand, strongly depends on the $\tan\beta$ value. While it is not a primary aim of the present paper to analyse such a scenario, when $\tan\beta \sim \mathcal{O}(10)$ the $\text{BR}(H^\pm \rightarrow \tau\nu_\tau)$ could be enhanced up to $\sim \mathcal{O}(0.01)$.

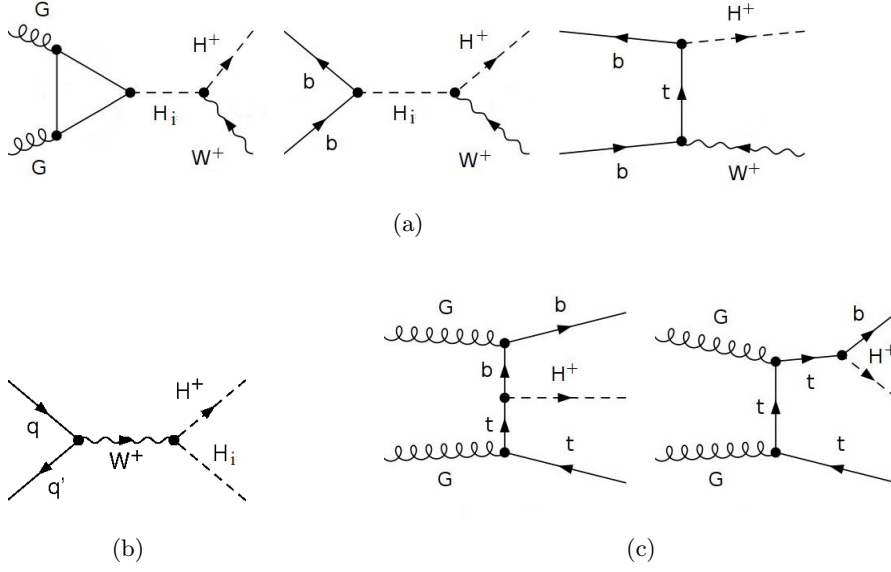


Figure 9: Single charged Higgs production channels at parton level.

$\sqrt{s} = 8$ or 14 TeV). Second, the intermediate W boson will be largely off-shell, also inducing significant depletion of the production rates. In contrast, the channel in fig. 9a receives contributions from both quark and gluon initiated processes and can further be resonant in the s -channel (for some of our benchmarks). In principle, the box contribution from heavy fermions should be included in the set, for it has been proven in several scenarios [119–121] that it can lead to $\mathcal{O}(10 - 100\%)$ corrections when $\tan\beta \sim \mathcal{O}(10)$. Still, in the phenomenological scenario that we propose to be tested at LHC ($\tan\beta \sim \mathcal{O}(1)$) this inclusion is totally irrelevant for our conclusions. The mode in fig. 9c also benefits from counting on two subchannels, though it is never resonant (as $M_{H^\pm} > m_t - m_b$ for the model considered here).

We consider now the four benchmark points P_2 , P_6 , P_7 and P_8 from table 1 and in fig. 10 we plot the cross sections against the charged Higgs mass for the aforementioned production channels. In the first place, we confirm that the associated production with a neutral scalar is disfavoured. Secondly, the remaining production mechanisms are always within a range of an order of magnitude at most. Again, we remark that the results are presented for an illustrative range of $M_H = 300 - 600$ GeV, while recalling that the allowed range is always restricted to the range $M_H \sim 380 - 450$ GeV for an intermediate choice of $\tan\beta$.

As regards the fermion-associated production mechanism of fig. 9c, we remark that it only depends on the values of M_{H^\pm} and $\tan\beta$ (see Eq. (2.13)), and there is a considerable reduction when moving from $\tan\beta = 1$ to $\tan\beta = 2$ (roughly a factor 2) due to the fact that the dominant contribution in the coupling is $\sim m_t / \tan\beta$, hence the ratio of VEVs acts as a reduction factor. The cross section of the fermion-associated contribution at $\tan\beta = 1$ is $\sim 10 - 10^2$ ($10^2 - 10^3$) fb when $\sqrt{s} = 8$ (14) TeV and it is mostly inversely proportional to $\tan\beta$.

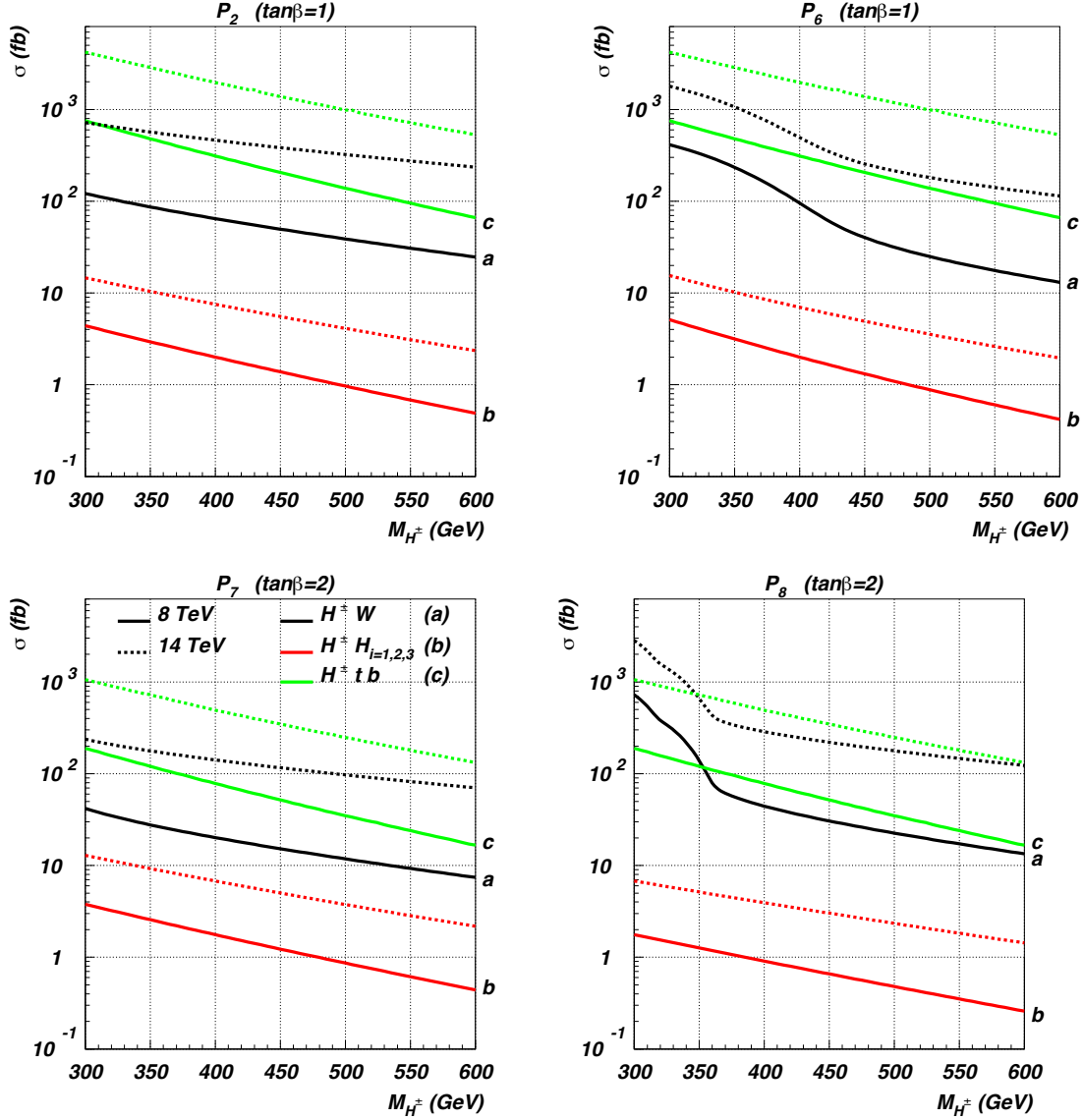


Figure 10: Cross sections of the single production mechanisms vs the charged Higgs mass at $\sqrt{s} = 8$ TeV (continuous lines) and $\sqrt{s} = 14$ TeV (dashed lines). Four benchmark points are considered, as indicated.

The scope of the fermion-associated production mechanism in extracting a $H^\pm \rightarrow W b \bar{b}$ signature (see below) has been analysed already in the literature, albeit in the MSSM, see [122], and we will revisit it in a CP-violating type-II 2HDM in a future publication.

Instead, here, we concentrate on vector-boson-associated production. The corresponding cross sections show a complicated behaviour with respect to different choices of parameters. We start our analysis by considering the channel with a final $H^\pm W^\mp$ state. From fig. 10 (P_2 and P_6) we see that a choice of $\tan\beta = 1$ plus a low-Higgs-masses scenario (P_2 : $M_2 \sim M_3 \sim 300 - 400$ GeV) has a cross section $\sim 10 - 10^2$ ($10^2 - 10^3$) fb when $\sqrt{s} = 8$ (14)

TeV, and that it is dominant (competitive) with respect to the fermion-associated production. On the other hand, we see that the choice of a high M_2 (P_6) and M_3 (≥ 500 GeV) favour the contribution from the parton-level channel $gg \rightarrow H_{i=2,3} \rightarrow H^\pm W^\mp$ proceeding through the on-shell $H_{i=2,3}$, and this results in a cross section that is always dominant and even enhanced when $m_W + M_{H^\pm} < M_{i=2,3}$, i.e. $\sim 10^2 - 10^3$ ($10^3 - 10^4$) fb when $\sqrt{s} = 8$ (14) TeV.

These qualitative conclusions also hold when $\tan \beta = 2$. Despite an overall suppression by one order of magnitude due to the increased value of $\tan \beta$, from fig. 10 (P_7) we see that on-shell production is not taking place (the H_2 and H_3 masses are too light). Hence, starting from a cross section of $\sim 10^2$ (10^3) fb when $\sqrt{s} = 8$ (14) TeV at the lower M_{H^\pm} scale we find a rate of ~ 10 (10^2) fb when $\sqrt{s} = 8$ (14) TeV for high values of M_{H^\pm} . However, the on-shell production is important for P_8 , where we can see an interplay between the on-shell H_2 and H_3 ($M_3 \sim 460$ GeV) production at $M_{H^\pm} \sim 300$ GeV being realised by a “double-shoulder” shaped line. In fact, when $M_{H^\pm} \sim 320$ (380) GeV the H_2 (H_3) on-shell production is switched off. In this framework, the cross section is ~ 10 (10^2) fb when $\sqrt{s} = 8$ (14) TeV, but it is increased by an order of magnitude when $M_{H^\pm} < 380$ GeV.

The high- $\tan \beta$ benchmark points give rather low production cross sections. For P_9 and P_{10} we find cross sections roughly half and a tenth, respectively, of that for P_8 . Since the latter, after the cuts discussed below, does not yield any useful signal, we have not explored P_9 and P_{10} any further.

Finally, we briefly comment on the sub-dominant channel with a W -mediated $H^\pm H_{i=1,3}$ final state. This channel is unlikely to have interesting phenomenological implication at the LHC, at least at the early running stage: when $\sqrt{s} = 8$ TeV and with integrated luminosity $\mathcal{L}_{\text{int}} = 10 \text{ fb}^{-1}$ the cross section is typically just above the threshold for producing a few events. Since it is not competitive with the other production channels, we will not study this mechanism.

In the next subsection we consider the neutral-Higgs-mediated production mechanism for analysing the scope of the LHC in discovering such a state in the allowed parameter space.

4.3 $pp \rightarrow H^\pm W^\mp$: significance analysis

In this subsection we analyse the significance of single charged Higgs boson production in association with gauge bosons for the set of benchmark points in table 1 (except P_9 and P_{10}). All figures herein refer to an integrated luminosity of 100 fb^{-1} .

Among the different charged Higgs decay modes, we have chosen to study $H^\pm \rightarrow W^\pm H_i$. The decay chain with the $H_i \rightarrow b\bar{b}$ intermediate decay is numerically favoured, so that we adopt it here, hence the complete H^\pm decay chain is

$$H^\pm \rightarrow W^\pm H_1 \rightarrow W^\pm b\bar{b}. \quad (4.4)$$

Therefore, we are interested in the significance ($\Sigma \sim S/\sqrt{B}$) analysis of a $2b+2W$ final state produced via a single charged Higgs. The most important background at the LHC for this final state is top quark pair production. However, we will show that a systematic reduction of this background is possible.

Notice that one may worry here about the contribution of the fermionic charged Higgs decay chain

$$H^+ \rightarrow t\bar{b} \rightarrow W^+ b\bar{b}, \quad (4.5)$$

as it yields an irreducible final state that is identical to the chosen one that could be defined as part of either the signal or the background. Under any circumstances, though, we believe that our top-mass veto (see below) will render this contribution negligible, so we omit it here³.

A b -tagging efficiency of $\sim 70\%$ has been assumed for each b -(anti)quark in the final state, and a full reconstruction efficiency has been assumed with respect to the W bosons. Among the possible decay patterns of the two W bosons, the semileptonic one was chosen, i.e., one hadronic and one leptonic decay, allowing for the full reconstruction of the events (unlike the fully leptonic decay mode) and a neater environment than the fully hadronic decay mode.

Hence, the overall selected process for the signal is the following:

$$pp \rightarrow W^\mp H^\pm \rightarrow W^\mp W^\pm H_1 \rightarrow W^\mp W^\pm b\bar{b} \rightarrow 2j + 2b + 1\ell + \text{MET}. \quad (4.6)$$

For each benchmark point, $2 \cdot 10^4$ unweighted events were produced. Regarding the top background, $4.5 \cdot 10^6$ unweighted events (with generation cuts) have been simulated in CalcHEP. For both signal and background the standard set of CTEQ6.6M [116] PDFs with scale $Q = \sqrt{s}$ were employed. For emulating a real LHC-prototype detector, a Gaussian smearing was included to take into account the electromagnetic energy resolution of $0.15/\sqrt{E}$ and the hadronic energy resolution of $0.5/\sqrt{E}$.

We describe now the overall strategy for the background reduction procedure. A first set of cuts includes typical detector kinematic acceptances and standard intermediate object reconstruction, such as $W \rightarrow jj$ and $H_1 \rightarrow b\bar{b}$ (cuts 1–3). Further, a t -(anti)quark reconstruction is used as “top veto” (cut 4). Led by the consideration that a b quark pair stemming from the Higgs boson is boosted (unlike the almost back-to-back pair from $t\bar{t}$), we define the last cut of the following set (cut 5):

- 1) **Kinematics:** standard detector cuts

$$\begin{aligned} p_\ell^T &> 15 \text{ GeV}, & |\eta_\ell| &< 2.5, \\ p_j^T &> 20 \text{ GeV}, & |\eta_j| &< 3, \\ |\Delta R_{jj}| &> 0.5, & |\Delta R_{\ell j}| &> 0.5; \end{aligned} \quad (4.7)$$

with η the pseudorapidity and $\Delta R = \sqrt{(\Delta\eta)^2 + (\Delta\phi)^2}$.

- 2) **light Higgs reconstruction:**

$$|M(b\bar{b}) - 125 \text{ GeV}| < 20 \text{ GeV}; \quad (4.8)$$

³We will attempt extracting this particular H^\pm topology in the context of the CP-violating 2HDM in a separate publication, with the aim of improving upon the MSSM results obtained in [123].

3) **hadronic W reconstruction ($W_h \rightarrow jj$):**

$$|M(jj) - 80 \text{ GeV}| < 20 \text{ GeV}; \quad (4.9)$$

4) **top veto:** if $\Delta R(b_1, W_h) < \Delta R(b_2, W_h)$, then

$$M(b_1 jj) > 200 \text{ GeV}, \quad M_T(b_2 \ell \nu) > 200 \text{ GeV}, \quad (4.10)$$

otherwise $1 \leftrightarrow 2$;

5) **same-hemisphere b quarks:**

$$\frac{\mathbf{p}_{b_1}}{|\mathbf{p}_{b_1}|} \cdot \frac{\mathbf{p}_{b_2}}{|\mathbf{p}_{b_2}|} > 0. \quad (4.11)$$

In table 3 we show the efficiency of the previous set of cuts against the simulated background for the P_2 and P_4 points of table 1, for two H^\pm masses. There is a clear correlation between the M_{H^\pm} value and the efficiency of the top veto (the most effective cut of this set): the higher the mass, the higher the efficiency.

Cut		BG events	BG Eff. (%)	P_2 events	P_2 Eff. (%)	S/\sqrt{B}
1:	Kin.	310291	100	54.2, 59.4	100, 100	0.1, 0.1
2:	H_1 rec.	263629	85.0	53.8, 59.0	99.3, 99.3	0.1, 0.1
3:	W rec.	256745	97.4	52.2, 57.4	97.0, 97.3	0.1, 0.1
4:	top veto	1689	0.7	15.5, 29.3	29.7, 51.0	0.4, 0.7
5:	same-side b 's	708	41.9	11.8, 22.8	76.1, 77.9	0.4, 0.9
Cut		BG events	BG Eff. (%)	P_4 events	P_4 Eff. (%)	S/\sqrt{B}
1:	Kin.	310291	100	356.2, 166.4	100, 100	0.6, 0.3
2:	H_1 rec.	263629	85.0	351.5, 165.7	98.7, 99.6	0.7, 0.3
3:	W rec.	256745	97.4	341.9, 160.4	97.3, 96.8	0.7, 0.3
4:	top veto	1689	0.7	41.6, 70.3	12.2, 43.8	1.0, 1.7
5:	same-side b 's	708	41.9	32.8, 54.7	78.7, 77.7	1.2, 2.1

Table 3: Consecutive efficiencies of the cuts imposed on the top quark background and (top) on the P_2 point, and (bottom) on the P_4 point (with $M_{H^\pm} = 310, 390 \text{ GeV}$) of table 1.

After these rather generic cuts are imposed, more signal-based selections can improve the significance. The main consideration of the following analysis is that the charged Higgs mass can equivalently be reconstructed by either the invariant mass of the four jets ($2b+2j$), $M(b\bar{b}jj)$, or the transverse mass of the b jets, the lepton and the MET, $M_T(b\bar{b}\ell\nu)$. Let's focus on the $M(b\bar{b}jj)-M_T(b\bar{b}\ell\nu)$ plane: for the signal, either of the two variables (if not both) will always reconstruct the correct charged Higgs boson mass, thus producing a cross-like shape in the plane defined by the two masses. In contrast, the background events

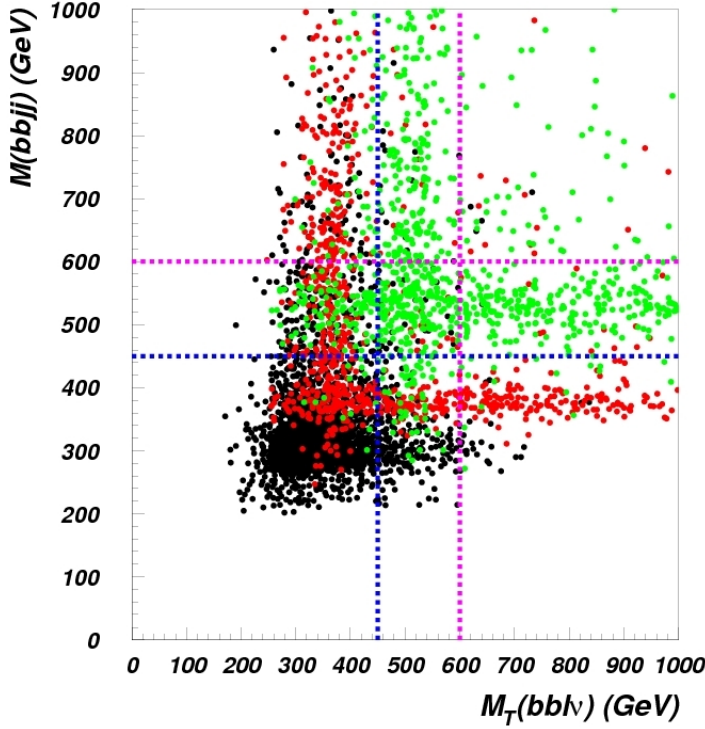


Figure 11: $M(b\bar{b}jj)$ vs. $M_T(b\bar{b}l\nu)$ after cut 5 for (unweighted) point P_2 , with $M_{H^\pm} = 375$ GeV (red, bottom-left) and $M_{H^\pm} = 525$ GeV (green, top-right). In black is the (unweighted) top background. The dashed lines show $M_{\text{lim}} = 450$ and 600 GeV.

accumulate at $\sim 2m_t$, as can be seen in fig. 11 in which we adopt an illustrative choice of charged Higgs masses.

The presence of long tails for the signal towards regions where the top background is heavily reduced allows us to introduce two specific (and alternative) cuts:

$$\text{“squared cut”}: \quad C_{\text{squ}} = \max(M(b\bar{b}jj), M_T(b\bar{b}l\nu)) > M_{\text{lim}} \quad (4.12)$$

$$\text{“single cut”}: \quad C_{\text{sng}} = M_T(b\bar{b}l\nu) > M_{\text{lim}}. \quad (4.13)$$

The single cut of eq. (4.13) is applied only on $M_T(b\bar{b}l\nu)$ because the reduction of the top background is higher than if compared to a similar cut on the $M(b\bar{b}jj)$ for the same numerical value of M_{lim} .

To determine which is the better of the two proposed strategies and what is the optimal value for M_{lim} , we studied the effects of C_{squ} and C_{sng} for several values of M_{lim} . Results are shown in tables 4 and 5 for the points P_2 and P_4 , respectively.

Clearly, a higher value for M_{lim} results in an increase of the significance, the top background is reduced more than the signal. It is important to note that for low charged Higgs masses, C_{squ} seems to perform better than the single cut. However, this is strictly true for $M_{H^\pm} \simeq 310$ GeV only: if a further selection is imposed, restricting the evaluation

Cut		$t\bar{t}$	$P_2 = 310 \text{ GeV}$		$P_2 = 390 \text{ GeV}$	
		Events	Events	S/\sqrt{B}	Events	S/\sqrt{B}
$M_{\text{lim}} = 450 \text{ GeV}$	C_{sng}	66.6	6.6	0.8	12.2	1.5
	C_{squ}	161.1	10.5	0.8	20.1	1.6
$M_{\text{lim}} = 500 \text{ GeV}$	C_{sng}	45.2	6.0	0.9	11.1	1.6
	C_{squ}	118.8	9.7	0.9	18.4	1.7
$M_{\text{lim}} = 550 \text{ GeV}$	C_{sng}	30.3	5.1	0.9	9.9	1.8
	C_{squ}	91.0	8.5	0.9	16.1	1.7
$M_{\text{lim}} = 600 \text{ GeV}$	C_{sng}	24.9	4.7	1.0	8.9	1.8
	C_{squ}	63.1	7.7	1.0	14.3	1.8

Table 4: Comparison between C_{squ} and C_{sng} vs M_{lim} for P_2 : surviving events and significance with respect to the background.

Cut		$t\bar{t}$	$P_4 = 310 \text{ GeV}$		$P_4 = 390 \text{ GeV}$	
		Events	Events	S/\sqrt{B}	Events	S/\sqrt{B}
$M_{\text{lim}} = 450 \text{ GeV}$	C_{sng}	66.6	14.5	1.8	29.0	3.6
	C_{squ}	161.1	25.8	2.0	47.3	3.7
$M_{\text{lim}} = 500 \text{ GeV}$	C_{sng}	45.2	12.7	1.9	26.3	3.9
	C_{squ}	118.8	22.4	2.1	43.0	3.9
$M_{\text{lim}} = 550 \text{ GeV}$	C_{sng}	30.3	10.8	2.0	23.4	4.2
	C_{squ}	91.0	19.8	2.1	37.9	4.0
$M_{\text{lim}} = 600 \text{ GeV}$	C_{sng}	24.9	10.0	2.0	20.3	4.1
	C_{squ}	63.1	17.7	2.2	33.1	4.2

Table 5: Comparison between C_{sng} and C_{squ} vs M_{lim} for P_4 : surviving events and significance with respect to the background.

of the significance to the peak-region only

$$\text{peak cut: } |M - M_{H^\pm}| < 50 \text{ GeV}, \quad (4.14)$$

the significance obtained by imposing C_{sng} , when calculated for all the other charged Higgs boson mass values, is *always higher* than the one obtained by imposing C_{squ} . Here, $M = \min(M(b\bar{b}jj), M_T(b\bar{b}\ell\nu))$ when eq. (4.12) is employed, while $M = M(b\bar{b}jj)$ when eq. (4.13) is employed.

For the following analysis, the value $M_{\text{lim}} = 600 \text{ GeV}$ has been chosen as well as the selection C_{sng} , this choice provides the best significance and a narrower peak while keeping a sufficient number of signal events (> 10). Should the surviving signal events be less than 10, it would then be advisable to choose instead the squared cut C_{squ} for the higher survival probability of the signal events (despite the lower significance and the broader peak).

The invariant mass distributions for the points P_2 , P_3 , P_4 , P_5 , and P_7 are plotted in

Figs. 12–16, each for two values of the charged Higgs mass. Table 6 collects the results for the S/\sqrt{B} analysis for all points.

	$M_{H^\pm} = 310 \text{ GeV}$		$M_{H^\pm} = 390 \text{ GeV}$	
	Events	S/\sqrt{B}	Events	S/\sqrt{B}
$t\bar{t}$	24.9			
peak	11.9	—	9.9	—
P_1	3.8	0.8	—	—
peak	2.6	0.8	—	—
P_2	4.7	1.0	8.8	1.8
peak	3.3	1.0	7.3	2.3
P_3	11.3	2.3	22.0	4.4
peak	7.7	2.3	17.2	5.4
P_4	10.0	2.0	20.3	4.1
peak	7.8	2.3	16.0	5.1
P_5	21.1	4.2	30.2	6.1
peak	13.9	4.1	25.0	7.9
P_6	14.0	2.8	—	—
peak	9.4	2.8	—	—
P_7	3.1	0.6	7.4	1.5
peak	2.8	0.8	7.3	2.3
P_8	1.2	0.2	—	—
peak	1.2	0.4	—	—

Table 6: Surviving events and their significance after the single cut of eq. (4.13) and after the peak selection of eq. (4.14), for all points of table 1, except P_9 and P_{10} .

As regards P_1 , P_6 and P_8 , no choice of allowed M_{H^\pm} produces any appreciable signal after the whole set of cuts, hence we will not discuss them any further.

From fig. 12 we learn that a choice of $\tan\beta = 1$ (though a rather low value of M_2 disfavors the production cross section) is enough to produce a visible signal, even for $M_{H^\pm} = 310 \text{ GeV}$ with the signal peak lying over the background. However, it is clear that the signal suffers from the selection cuts, and this critical situation is eased up only when $M_{H^\pm} = 390 \text{ GeV}$. In fact, for higher allowed M_{H^\pm} masses, the signal contains a conspicuous number of events (~ 10 in the peak bin), and it is always above the background.

If we consider the points P_3 , P_4 and P_5 the same circumstances occur: the higher M_{H^\pm} , the more visible and clear the signal with respect to the background. Moreover, from P_3 ($M_2 = 350 \text{ GeV}$) to P_5 ($M_2 = 450 \text{ GeV}$) the production cross section increases, as shown in figs. 13–15.

If we consider the point P_7 (which has $\tan\beta = 2$) in fig. 16, we do not note any change from the previous considerations: even when $\tan\beta$ grows we still have a considerable production cross section in the allowed parameter space, and the signal would be observable with respect to the background, at least for $M_{H^\pm} = 390 \text{ GeV}$.

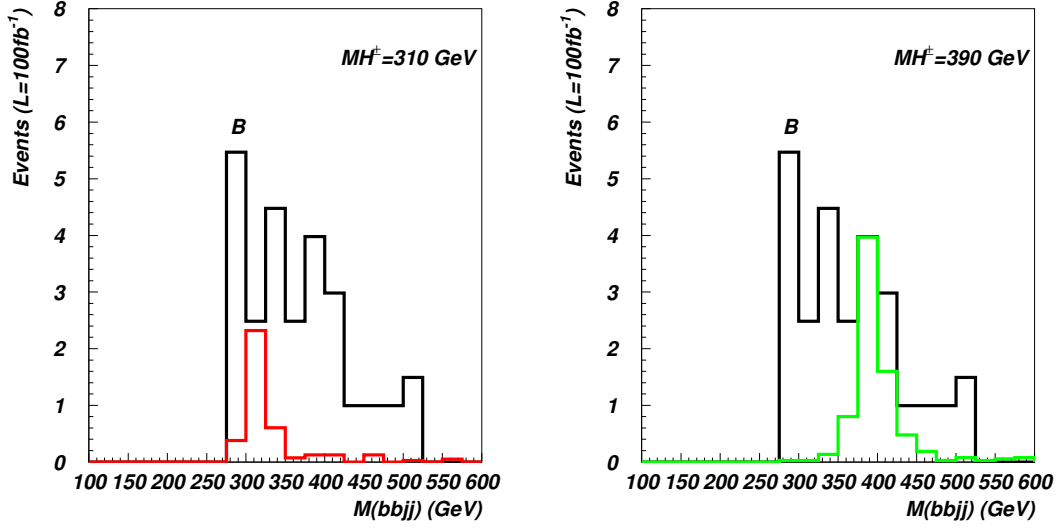


Figure 12: Point P_2 . Number of events integrated with $\mathcal{L}_{\text{int}} = 100 \text{ fb}^{-1}$ at $\sqrt{s} = 14 \text{ TeV}$ vs $M(b\bar{b}jj)$ for signal (coloured lines) and t -quark background (black histogram, labeled “B”).

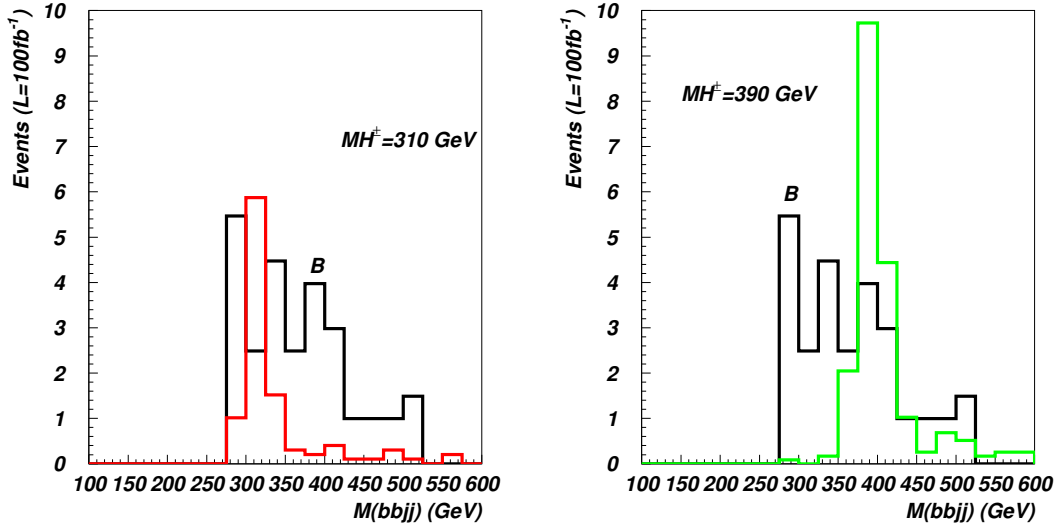


Figure 13: Point P_3 . Similar to fig. 12.

5 Possible future scenarios

We shall here discuss possible future experimental developments, and consider their implications for the model, in particular for the proposed benchmarks. The basic question is of course: which experimental efforts are required to exclude the 2HDM altogether? We shall below adress a couple of LHC-related aspects of this question.

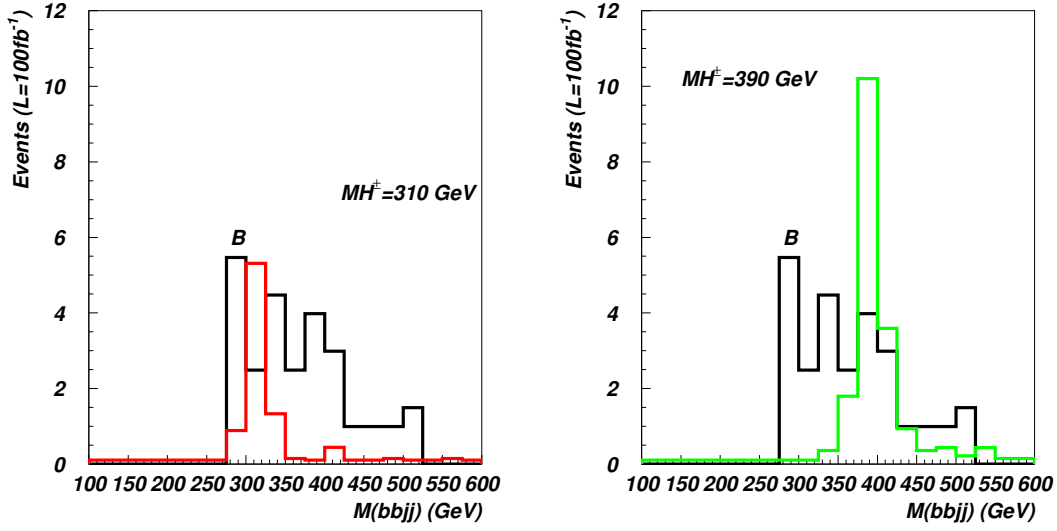


Figure 14: Point P_4 . Similar to fig. 12.

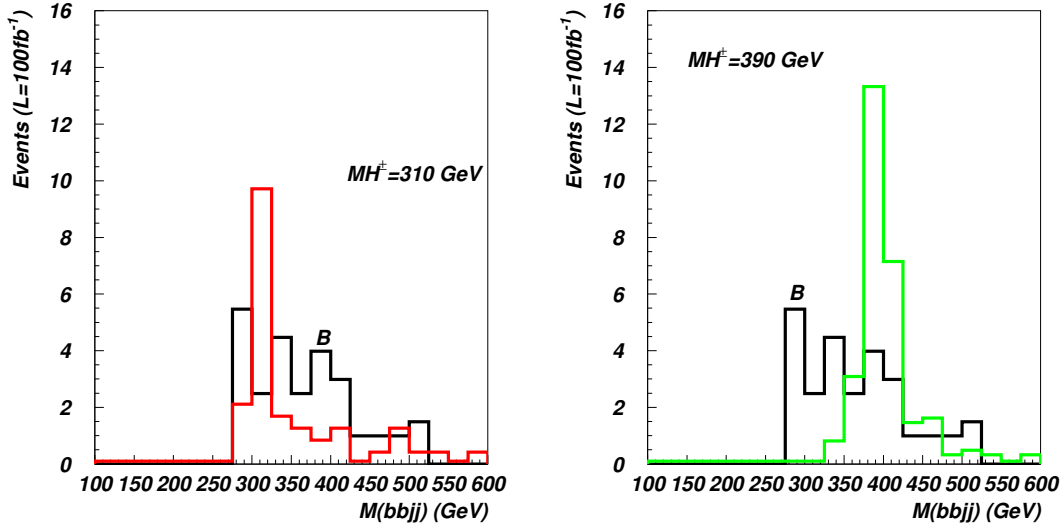


Figure 15: Point P_5 . Similar to fig. 12.

5.1 Higher and more constrained rates for $gg \rightarrow H_1 \rightarrow \gamma\gamma$

Several authors have recently argued that the LHC experiments point to an overall rate for $pp \rightarrow H \rightarrow \gamma\gamma$ that is somewhat high compared to the SM prediction. In our parameter scans in section 3 we generously allowed the ratio $R_{\gamma\gamma}$ of Eq. (3.18) to satisfy $0.5 \leq R_{\gamma\gamma} \leq 2.0$. We shall here briefly comment on how the parameter space is further constrained when we only allow the upper range:

$$1.5 \leq R_{\gamma\gamma} \leq 2.0 \quad (5.1)$$

For the case shown in Fig. 2, namely $\tan\beta = 1$ and $M_2 = 400$ GeV (and two values of

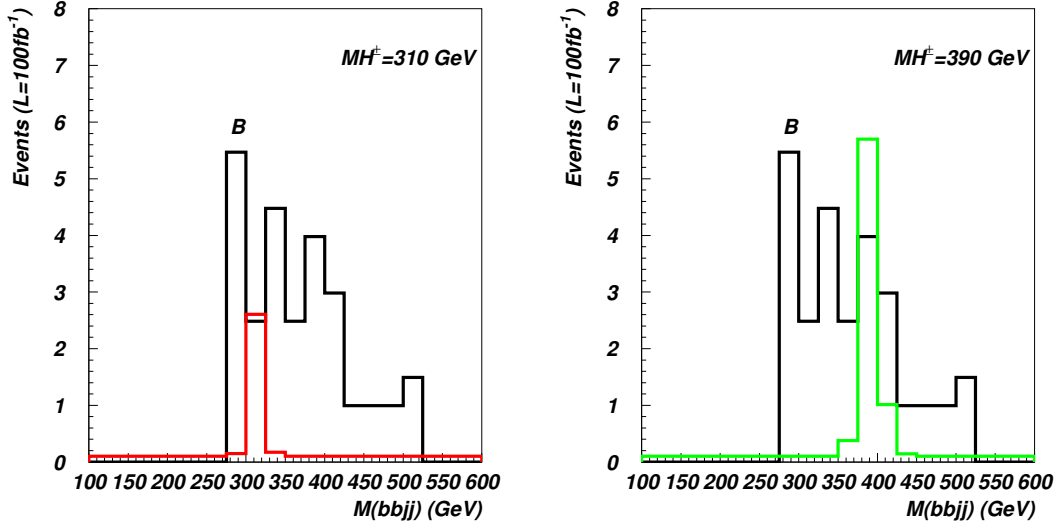


Figure 16: Point P_7 . Similar to fig. 12.

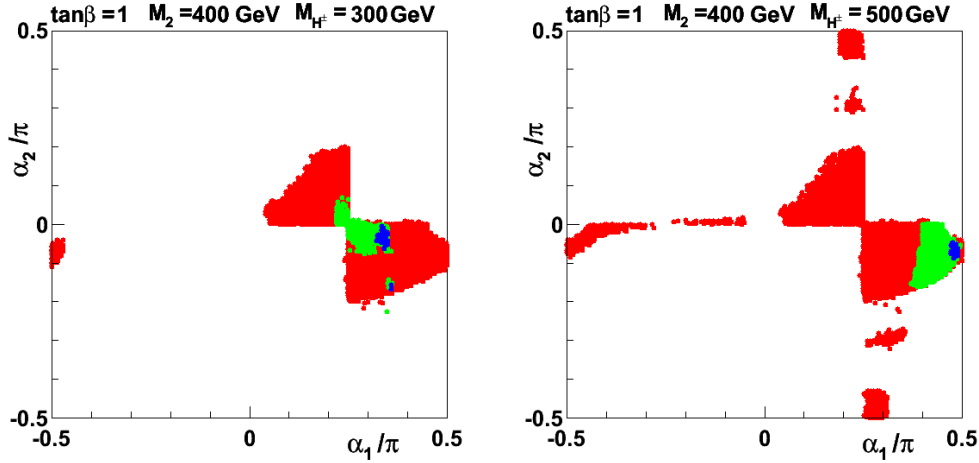


Figure 17: Superimposed on Fig. 2, we show (in blue) the remaining allowed regions in the α_1 - α_2 parameter space, imposing higher rates for the $\gamma\gamma$ channel, Eq. (5.1), for $M_2 = 400$ GeV and the additional parameters given in Eq. (3.26).

M_{H^\pm}), we show in Fig. 17 how the LHC-allowed region gets constrained. Such a development can have dramatic consequences for the model: benchmark points P_1 , P_6 , P_7 and P_8 would be excluded.

In figure 18 we show the remaining allowed regions in the $\tan\beta$ - M_{H^\pm} plane, which exhibits a significant reduction, as compared with figure 6. While the scans have limited statistics, there is an indication that the remaining allowed parameter space starts fragmenting into disconnected regions.

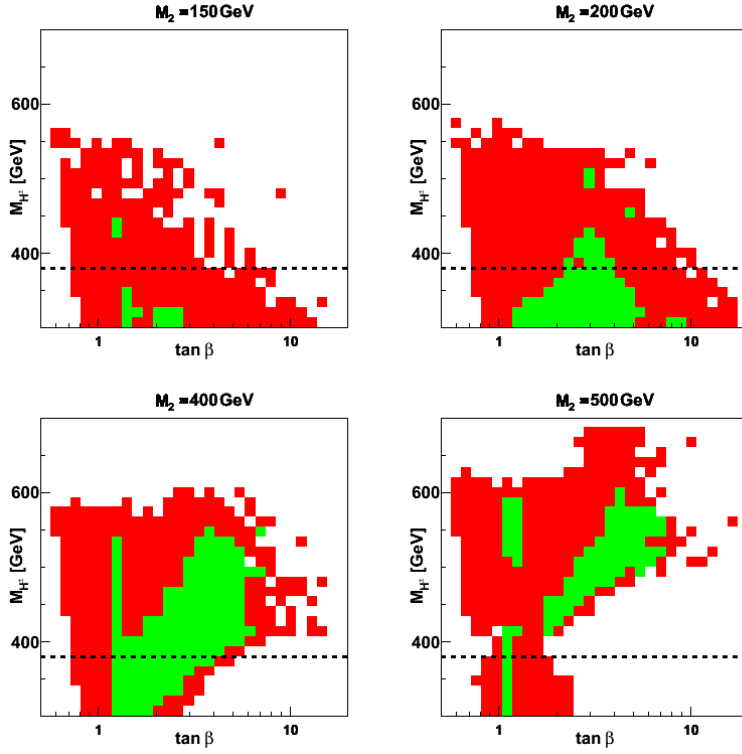


Figure 18: Allowed regions in the $\tan\beta$ – M_{H^\pm} parameter space, without (red) and with (green) the LHC constraints, for the range in $R_{\gamma\gamma}$ given by Eq. (5.1), and four values of M_2 , as indicated. The dashed lines show the recent bound at 380 GeV [69].

5.2 Tightened upper bound on $gg \rightarrow H_{2,3} \rightarrow W^+W^-$ (and ZZ)

Presumably, the search for an SM-like Higgs will continue in the mass range from around 130 GeV and up. To a first (rough) approximation, a Higgs in this mass region is produced via gluon fusion, and decays via WW (or ZZ) bosons. Assuming these upper bounds are tightened, it is interesting to see how the allowed parameter space behaves. In Fig. 19 we show how the allowed parameter space shrinks if we assume that the upper bound on a Higgs-like particle, represented by the quantity R_{ZZ} of Eq. (3.20), is lowered by a factor 0.5. Again, we see a rather dramatic impact of such a development.

6 Conclusions

For the channel $pp \rightarrow H^\pm W^\mp \rightarrow W^+W^-b\bar{b}$, we have established a set of 7 benchmarks for the CP-violating 2HDM with type-II Yukawa interactions. These points all have $M_1 = 125$ GeV, low $\tan\beta$, they all violate CP, and allow for a range of charged-Higgs masses. A set of cuts is proposed, that will reduce the $t\bar{t}$ background to a tolerable level, and allow for the detection of a signal in the $WW \rightarrow jj\ell\nu$ channel. Some of the proposed benchmark points lead to enhanced $H^\pm W^\mp$ production cross sections due to resonant production via H_2 or H_3 in the s -channel.

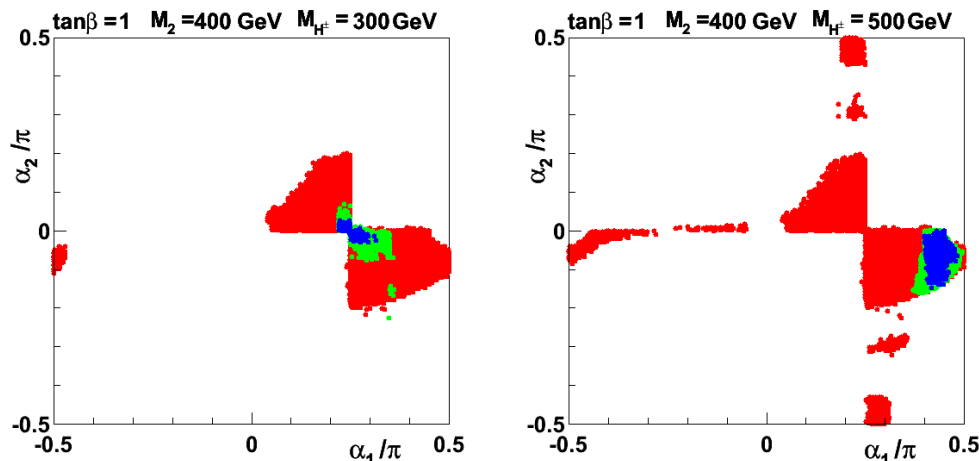


Figure 19: Superimposed on Fig. 2, we show (in blue) the remaining allowed regions in the α_1 – α_2 parameter space, assuming that the upper bound on an SM-like Higgs is tightened by a factor 0.5 for masses in the range 130–600 GeV, for $M_2 = 400$ GeV and the additional parameters given in Eq. (3.26).

Most of the proposed points are in the interior of some allowed domain in the α space, and thus robust with respect to minor modifications of the experimental constraints. Some of the benchmark points are vulnerable to a higher value of $R_{\gamma\gamma}$. However, the points P_2 , P_3 , P_4 and P_5 are not endangered.

It should also be noted that the proposed channel only benefits from favourable production cross sections and branching ratio at low values of $\tan\beta$. In this region, the charged Higgs mass is constrained to the range ~ 380 – 470 GeV.

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A Appendix. The decoupling limit

We shall here explore the so-called decoupling limit, which has been studied for the CP-conserving case in [111], where

$$M_{H^\pm} \sim M_3 \sim M_2 \gg M_1. \quad (\text{A.1})$$

We shall see that the large masses will all be of order μ . Note, however, that this discussion disregards the constraints of positivity, unitarity, etc., that are discussed in section 3.

A.1 $\tan \beta$ of $\mathcal{O}(1)$

For definiteness, we substitute

$$M_3 = M_2 = M \quad (\text{A.2})$$

into the expressions for the λ 's [34], require them all to be small, and also neglect terms of order M_1^2 compared to M^2 . The conditions related to the different λ 's are:

$$\lambda_1 : \quad (c_1^2 s_2^2 + s_1^2) M^2 \simeq s_\beta^2 \mu^2, \quad (\text{A.3a})$$

$$\lambda_2 : \quad (s_1^2 s_2^2 + c_1^2) M^2 \simeq c_\beta^2 \mu^2, \quad (\text{A.3b})$$

$$\lambda_3 : \quad \frac{c_1 s_1}{c_\beta s_\beta} c_2^2 M^2 \simeq 2M_{H^\pm}^2 - \mu^2, \quad (\text{A.3c})$$

$$\lambda_4 : \quad c_2^2 M^2 \simeq 2M_{H^\pm}^2 - \mu^2, \quad (\text{A.3d})$$

$$\text{Re } \lambda_5 : \quad c_2^2 M^2 \simeq \mu^2, \quad (\text{A.3e})$$

$$\text{Im } \lambda_5 : \quad (c_\beta c_1 + s_\beta s_1) c_2 s_2 \simeq 0. \quad (\text{A.3f})$$

Adding Eqs. (A.3a), (A.3b) and (A.3e), we get

$$M^2 \simeq \mu^2, \quad (\text{A.4})$$

as anticipated. Substituting back into Eqs. (A.3e) and (A.3d), we find

$$c_2 \simeq 1, \quad (\text{A.5})$$

and

$$M_{H^\pm}^2 \simeq \mu^2 \simeq M^2. \quad (\text{A.6})$$

The last equation, Eq. (A.3f), provides two solutions, either

$$\cos(\beta - \alpha_1) \simeq 0, \quad \text{or} \quad \sin(2\alpha_2) \simeq 0. \quad (\text{A.7})$$

We note that the angle α_3 does not enter in these asymptotic conditions (A.3), and that they are all satisfied for

$$\alpha_1 \sim \beta, \quad \alpha_2 \sim 0, \quad \alpha_3 \text{ arbitrary}. \quad (\text{A.8})$$

A.2 Large $\tan\beta$

The case of large values of $\tan\beta$ requires special attention. Because of over-all factors which were left out in Eq. (A.3), the first on them, Eq. (A.3a) must be satisfied to a much higher degree than the others (a factor $1/c_\beta^2$ is involved). This means that the expression

$$c_1^2 s_2^2 + s_1^2 \quad (\text{A.9})$$

must be maximised. This requires

$$\sin\alpha_1 = \pm 1, \quad \sin\alpha_2 = 0, \quad (\text{A.10})$$

consistent with Eq. (A.8) (and with H_2 or H_3 being odd under CP), or

$$\sin\alpha_1 = 0, \quad \sin\alpha_2 = \pm 1. \quad (\text{A.11})$$

The latter solution (which corresponds to H_1 being odd under CP) is not contained in Eq. (A.8).

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